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SOVIET PHYSICS

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DETERMINATION OF THE ELECTRON AFFINITY OF SULFUR BY SURFACE IONIZATION

I. N. Bakulina and N. I. Ionov

(Presented by Academician A. A. Lebedev April 27, 1957)

In an earlier paper [1] we described experiments with the surface ionization of potassium thiocyanate (KCNS) on wolfram, in which we observed the formation of the negative ions CN^- and S^-. A very rough estimate of the electron affinity of sulfur gave a value of approximately 2 ev. In the present paper we describe experiments which provide a more accurate value.

Our method of determining the electron affinity was used in previous work [2]. It is based on the comparison of negative ion currents of two elements which are ionized simultaneously on the same surface. The current I of negative ions from an ionizing surface is known to be given by the formula

$$I = \frac{n\varepsilon}{1 + \frac{g_a}{g_p} e^{\frac{\varepsilon(\varphi - s)}{kT}}}, \qquad (1)$$

where n is the number of atoms reaching the surface in unit time; T is the temperature of the surface; s is the electron affinity of the atoms; φ is the work function of the surface; g_p and g_a are the statistical weights of the element in the ionic and atomic states, respectively; k is the Boltzmann constant; ε is the electron charge. We investigated ionization on the surface of a polycrystalline wolfram filament, for which the work function of individual parts can vary within wide limits between φ_{min} and φ_{max}. However, the electron affinity of even the most electronegative atoms, such as the halogens, is in the case of pure wolfram so much smaller than φ_{min} that for all possible temperatures of the wolfram it is possible in the denominator of Eq. (1) to neglect unity compared with the exponential terms. Equation (1) then becomes

$$I \cong n\varepsilon \frac{g_p}{g_a} e^{\frac{\varepsilon(s-\varphi)}{kT}}. \qquad (2)$$

If in unit time n_1 atoms of one element and n_2 atoms of another element reach the same surface, by measuring the current ratio I_1/I_2 of the negative ions from the surface and the surface temperature T we can determine the difference of the electron affinities of the ionizing atoms by the formula

$$\frac{I_1}{I_2} = \frac{n_1}{n_2} \quad \frac{A_2}{A_1} e^{\frac{\varepsilon(s_1 - s_2)}{kT}}, \tag{3}$$

whence

$$s_1 - s_2 = \frac{T}{5040}\left[\log \frac{I_1}{I_2} + \log \frac{n_2 A_2}{n_1 A_2}\right], \tag{4}$$

where $A_1 = g_{a1}/g_{p1}$, $A_2 = g_{a2}/g_{p2}$. Thus from a knowledge of the electron affinity of one element we can determine the electron affinity of the other.

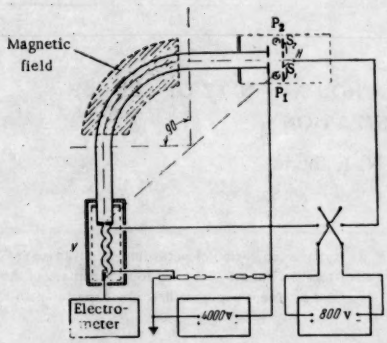

Fig. 1.

In this work we determined the difference between the electron affinities of bromine and sulfur by studying the surface ionization of NaBr and Na_2S. As previously [2], molecular beams of NaBr and Na_2S from two independent platinum furnaces P_1 and P_2 (Fig. 1) were directed against the surface of a wolfram filament H heated to temperature T. The positive and negative ions formed on the filament were analyzed by a sector-type magnetic mass spectrometer. The molecular beams could be stopped independently by means of the shutters S_1 and S_2.

The mass spectrometric analysis showed that when NaBr and Na_2S were ionized the positive ion spectrum contained only Na^+ ions while the negative ion spectrum contained Br^- and S^-. For equal densities of the NaBr and Na_2S molecular beams the current of negative sulfur ions at all filament temperatures was smaller by two to three orders of magnitude than the current of negative bromine ions and was $\sim 10^{-15}$ amp in absolute magnitude. For the measurement of such small currents we used a box-type ion-electron multiplier (Y) with open entrance, the first dynode of which was reached by the ion beams after emerging from the mass analyzer. The dynodes were made of a Cu-Be alloy and were only slightly activated by thermal treatment. The total multiplication factor of the multiplier in 12 steps was ~ 2000, which did not change with time and was very stable against the effects of the air. The electron current from the collector plate was fed to the input of an ordinary EMU-3 electrometric amplifier. The minimum ion current which could be recorded reliably with this instrument was $\sim 10^{-17}$ amp.

The experimental procedure was as follows. Pure nonaqueous Na_2S was prepared from the original crystal hydrate $Na_2S \cdot 9H_2O$ [3]. The Na_2S and chemically pure NaBr were placed in separate platinum furnaces and gradually heated in a high vacuum to higher than the working temperatures. Before measurement of the Br^- and S^- currents the positive sodium ion currents from alternate ionization of the NaBr and Na_2S beams were equalized. Additivity of the Na^+ currents from both beams was also tested (with both shutters open).

We assume that at wolfram temperatures above 1800°K the Na_2S and NaBr are completely dissociated and equality of the positive sodium currents therefore means that in each second there are formed on the wolfram surface n_1 free atoms of bromine and $n_2 = {}^1/_2\, n_1$ sulfur atoms, which are partially ionized in accordance with (1). For all measurements of the current ratio I_1/I_2 (where I_1 is the Br^- current and I_2 is the S^- current) a periodic check was made of the constancy of molecular beam densities from the furnaces through measurement of the corresponding Na^+ currents.

The temperature of the wolfram filament was determined by an optical pyrometer. From measurements of the I_1/I_2 ratio at various filament temperatures T we calculated by means of (4) the difference between the electron affinities of bromine and sulfur. Equation (4) also includes the statistical weights of bromine and sulfur in the atomic and ionic states. In the ground state the statistical weight of the bromine atom is $g_{a1} = 4$; that of the negative bromine ion is $g_{p1} = 1$; for the sulfur atom $g_{a2} = 5$ and for the sulfur ion $g_{p2} = 4$. But the sulfur atom has two excited states $^3P_1$ and $^3P_0$ close to the ground state, allowance for which in the temperature range from

1800 to 2300°K gives the statistical weight $g_{a1} = 8$. Excited states of the bromine atom up to 2300°K can be neglected and for the negative Br^- and S^- ions they are unknown. Therefore in calculation $s_1 - s_2$ from (4) we assumed $A_1/A_2 = 2$, allowing for excited states of the sulfur atom.

From the experimental values of I_1/I_2 we calculated $s_1 - s_2$ for various temperatures of the wolfram filament from 1800 to 2300°K. The average value obtained from more than 80 independent measurements was

$$s_1 - s_2 = 1.23 \pm 0{,}05 \text{ ev}.$$

The mean arithmetic error of all determinations is given. This agrees with our estimate of the accuracy of measuring I_1/I_2 and the filament temperature.

However, possible differences in the secondary ion-electronic emission of Br^- and S^- at the first dynode of the multiplier and in the transmission of the mass spectrometer for these ions could result in a systematic error of the I_1/I_2 measurements. A special experiment showed that these parameters agree to within $\pm 20\%$ so that the systematic error in determination of $s_1 - s_2$ does not exceed 0.03 ev and the over-all error is less than 0.1 ev.

Our method does not yield absolute values of the electron affinity. Therefore, if we assume as previously [2] that the electron affinity of bromine atoms is 3.6 ev the electron affinity of sulfur atoms according to our measurements is $s_2 = 2.37$ ev. The accuracy of this value depends not only on the error of $s_1 - s_2$ but also on the difference between the true value of s_1 for bromine atoms and our assumed value.

Our value for the electron affinity of sulfur atoms agrees with the value $s_2 > 2.2$ ev obtained in experiments on the ionization of SO_2 molecules by electron impact [4]. However, recently performed determinations of the electron affinity of sulfur by electron photodetachment from S^- ions gave $s_2 = 2.07 \pm 0.07$ ev [5]. The difference from our value may possibly indicate that the electron affinity of bromine is below 3.6 ev, provided, of course, that both methods do not conceal some additional systematic errors.

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RADIATION DIFFUSION IN A SEMI-INFINITE MEDIUM

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(Presented by Academician V. A. Ambartsumian April 16, 1957)

The problem of radiation diffusion in a semi-infinite medium is encountered in astrophysics (in connection with the study of stellar and planetary atmospheres) and in geophysics (in the study of water basins). The different special cases are usually considered separately. In the present note we shall show that the solutions of all problems of radiation diffusion in semi-inifinite media differing in the arrangement of the radiation sources can be expressed by the same function, depending only on the optical depth.

We shall assume that isotropic radiation scattering occurs in a medium, with the probability λ that a quantum will survive an elementary scattering. The radiation field in this medium is then obtained by solving the following integral equation for the function $B(\tau)$:

$$B(\tau) = \frac{\lambda}{2}\int_0^\infty B(\tau')\,\mathrm{Ei}\,|\tau - \tau'|\,d\tau' + g(\tau), \tag{1}$$

where $g(\tau)$ gives the arrangement of the sources. When $B(\tau)$ is found the intensity of the radiation at optical depth τ and forming the angle ϑ with the normal is expressed by

$$I(\tau, \vartheta) = \int_\tau^\infty B(\tau')\, e^{-(\tau'-\tau)\sec\vartheta}\sec\vartheta\, d\tau' \quad \left(\vartheta < \frac{\pi}{2}\right); \tag{2}$$

$$I(\tau, \vartheta) = -\int_0^\tau B(\tau')e^{-(\tau'-\tau)\sec\vartheta}\sec\vartheta\, d\tau' \quad \left(\vartheta > \frac{\pi}{2}\right). \tag{3}$$

The formal solution of (1) is

$$B(\tau) = g(\tau) + \int_0^\infty \Gamma(\tau, \tau')\, g(\tau')\, d\tau', \tag{4}$$

where $\Gamma(\tau, \tau')$ is the resolvent. Using the probability significance of the resolvent it is possible to obtain ([1], p. 225) the following equation for its determination:

$$\frac{\partial\Gamma}{\partial\tau} + \frac{\partial\Gamma}{\partial\tau'} = \Gamma(\tau, 0)\,\Gamma(0, \tau'), \tag{5}$$

which gives (for $\tau' > \tau$):

$$\Gamma(\tau, \tau') = \Phi(\tau' - \tau) + \int_0^\tau \Phi(x)\,\Phi(x + \tau' - \tau)\, dx, \tag{6}$$

where $\Gamma(\tau, 0) = \Phi(\tau)$. The definition of $\Phi(\tau)$ is given by the integral equation

$$\Phi(\tau) = K(\tau) + \int_0^{\tau} K(\tau - \tau')\,\Phi(\tau')\,d\tau', \tag{7}$$

where

$$K(\tau) = \frac{\lambda}{2}\int_0^1 e^{-\tau/\eta}\,\varphi(\eta)\,\frac{d\eta}{\eta}, \tag{8}$$

and $\varphi(\eta)$ is Ambartsumian's function, which he defined by [2]

$$\varphi(\eta) = 1 + \frac{\lambda}{2}\eta\varphi(\eta)\int_0^1 \varphi(\zeta)\,\frac{p\zeta}{\eta+\zeta}. \tag{9}$$

The Laplace transform of (7) gives

$$\int_0^{\infty} \Phi(\tau)\,e^{-m\tau}\,d\tau = \frac{1}{1 - \frac{\lambda}{2}\int_0^1 \varphi(\eta)\,\frac{d\eta}{1+m\eta}} - 1. \tag{10}$$

From (10) with $m = \frac{1}{\zeta} \geq 1$ it follows that

$$\int_0^{\infty} \Phi(\tau)\,e^{-\tau/\zeta}\,d\tau = \varphi(\zeta) - 1. \tag{11}$$

Inversion of the Laplace transformation leads to the following asymptotic formula for $\Phi(\tau)$ when τ is large:

$$\Phi(\tau) = \frac{e^{-k\tau}}{\frac{\lambda}{2}\int_0^1 \varphi(\eta)\,\frac{\eta d\eta}{(1-k\eta)^2}}, \tag{12}$$

where k is related to λ by

$$\frac{\lambda}{2k}\lg\frac{1+k}{1-k} = 1. \tag{13}$$

From (10) it is also possible to obtain both exact and approximate formulas for $\Phi(\tau)$, as will be discussed later.

Thus the determination of the radiation field in a semi-infinite medium with arbitrary sources of radiation is reduced to the finding of $\Phi(\tau)$. When this function is known (6) gives the resolvent $\Gamma(\tau, \tau')$, Eq. (4) gives $B(\tau)$ and (2) and (3) give the radiation intensity $I(\tau, \vartheta)$.

Let us consider some special cases of radiation diffusion in a semi-infinite medium.

1) Let

$$g(\tau) = Ge^{-m\tau}, \tag{14}$$

where G and m are constants. Using (4), (5) anf (10), we obtain

$$B(\tau, m) = \frac{G}{1 - \frac{\lambda}{2}\int_0^1 \varphi(\eta)\,\frac{d\eta}{1+m\eta}}\left[e^{-m\tau} + \int_0^{\tau} \Phi(\tau')\,e^{-m(\tau-\tau')}\,d\tau'\right] \tag{15}$$

Equation (15) is valid for both positive m and negative values with $-m < k$.

When the medium is illuminated by parallel rays forming the angle arc cos ζ with the normal and producing the illumination πS on an area at right angles $g(\tau) = \frac{\lambda}{4} S e^{-\tau/\zeta}$. Setting in (15) $G = \frac{\lambda}{4} S$ and $m = \frac{1}{\zeta} \geq 1$, we obtain

$$B\left(\tau, \frac{1}{\zeta}\right) = \frac{\lambda}{4} S \varphi(\zeta) \left[e^{-\tau/\zeta} + \int_0^\tau \Phi(\tau') e^{-(\tau-\tau')/\zeta} d\tau' \right] \tag{16}$$

which agrees with earlier results ([1], p. 112).

2) Let

$$g(\tau) = \tau^n, \tag{17}$$

where n is a positive integer. Using (4) and (5), we obtain the recurrence relation

$$B_n(\tau) = n \int_0^\tau B_{n-1}(\tau') d\tau' + \Psi(\tau) \int_0^\infty \Phi(\tau') \tau'^n d\tau', \tag{18}$$

where

$$\Psi(\tau) = 1 + \int_0^\tau \Phi(\tau) d\tau. \tag{19}$$

The last integral in (18) is easily expressed by means of (7) and (8) in terms of the moments of $\varphi(\eta)$,

$$\alpha_n = \int_0^1 \varphi(\eta_1) \eta_1^n d\eta_1. \tag{20}$$

For a uniform distribution ($g(\tau) = 1$) we have

$$B_0(\tau) = \frac{\Psi(\tau)}{\sqrt{1-\lambda}}, \tag{21}$$

and when $g(\tau) = \tau$

$$B_1(\tau) = \frac{1}{\sqrt{1-\lambda}} \int_0^\tau \Psi(\tau') d\tau' + \frac{\lambda}{2} \frac{\alpha_1}{1-\lambda} \Psi(\tau) \tag{22}$$

etc.

3) Let us assume that pure radiation scattering ($\lambda = 1$) takes place in the medium and that the sources are located at infinite depth. Denoting $B(\tau)$ for this case by $B_*(\tau)$ and taking into account the probability significance of $\Gamma(\tau, \tau')$, we obtain $B_*(\tau) \sim \Gamma(\tau, \infty)$, or by using (6),

$$B_*(\tau) = B_*(0) \Psi(\tau). \tag{23}$$

Equation (23) is, specifically, the solution for radiation equilibrium of a stellar photosphere. With πF denoting the radiation flux in the photosphere (independent of τ) we have

$$F = 2 \int_0^\infty \mathrm{Ei}_2 \tau B_*(\tau) d\tau = 2 B_*(0) \int_0^\infty \mathrm{Ei}_2 \tau \Psi(\tau) d\tau. \tag{24}$$

Substituting (19) and (11) we obtain the familiar relation

$$F = 2B_*(0)\,\alpha_1 = \frac{4}{\sqrt{3}} B_*(0). \tag{25}$$

In the theory of photospheres $B_*(\tau)$ is usually represented by

$$B_*(\tau) = \frac{3}{4} F[\tau + q(\tau)]. \tag{26}$$

Equating (23) and (26), we arrive at the following relation between $q(\tau)$ and $\Psi(\tau)$:

$$q(\tau) = \frac{\Psi(\tau)}{\sqrt{3}} - \tau. \tag{27}$$

Thus in all of the problems considered here $B(\tau)$ is expressed very simply in terms of $\Phi(\tau)$ or $\Psi(\tau)$.

When we know either the exact or approximate solution of any single problem and obtain the expression of $\Phi(\tau)$ from it, we can use the foregoing formulas to obtain solutions of other problems. As an example let us take the approximate solution of $B_*(\tau)$ obtained by Chandrasekhar [3]:

$$B_*(\tau) \simeq \frac{3}{4} F\left(\tau + \frac{1}{\sqrt{3}}\right). \tag{28}$$

It is easily seen that in this approximation

$$\Phi(\tau) \simeq \sqrt{3}, \quad \varphi(\eta) \simeq 1 + \eta\sqrt{3}. \tag{29}$$

Inserting (29) into (16), we obtain an approximate solution for the diffusion of radiation in a medium illuminated by parallel rays ($\lambda = 1$):

$$B\left(\tau, \frac{1}{\zeta}\right) \simeq \frac{S}{4}(1 + \zeta\sqrt{3})\,[e^{-\tau/\zeta} + \sqrt{3}(1 - e^{-\tau/\zeta})\zeta]. \tag{30}$$

The results obtained here can easily be extended to a medium of finite optical thickness.

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VITRIFICATION OF FLUCTUATIONS AND THE SCATTERING OF LIGHT

M. V. Vol'kenshtein

(Presented by Academician A. F. Ioffe May 15, 1957)

The characteristics of molecular scattering of light in glasses become evident when the glass is compared with a liquid above the vitrification temperature T_g. Light scattering in silicate glasses is hardly molecular in the true sense, but the indicated comparison is difficult because of the high values of T_g. T. S. Velichkina [1] has studied light scattering by triacetin in β-chloroethyl ether of benzosulfo acid over a broad temperature range above and below T_g. The scattering intensities for glasses were somewhat lower than for liquids when $T > T_g$ (by 20 to 30%) but were 100 to 130% higher than the calculated values for glasses regarded as liquids according to Einstein's formula as corrected by Cabannes [2]. The experimental intensities were one order of magnitude greater than the theoretical values for isotropic solids [3]. For fused quartz the intensity was 8 times greater than the value calculated from Einstein's formula. However, when we insert into this formula not the experimental T but $T_g = 2100°K$ we obtain a value 20% below the experimental figure [4].

Contrary to the Landau-Placzek formula [5] the central component of the Rayleigh line of viscous liquids and glasses is very intense and narrow. Shifted components are practically absent in glasses [6-8]. This is also observed in substances with small molecular anisotropy, in which case the central component is polarized [9].

An explanation of these facts must be sought on the basis of kinetic concepts. Despite the superficial resemblance of vitrification and second-order phase transitions, the same theory of light scattering cannot be applied [10]. Vitrification of a liquid is a kinetic process [11]. Consideration of a relaxation process characterized by the time τ gives the following condition for vitrification:

$$\left(\frac{d\tau}{dT}\right)_{T=T_g} = \frac{1}{q}. \tag{1}$$

where $q = -|dT/dt|$ is the cooling rate [12].

The peculiarities of light scattering in glasses are explained by the vitrification of the fluctuations of structural entropy. In a liquid consisting of isotropic molecules light scattering is determined by the fluctuations of the specific volume which can be represented by

$$\Delta v = \Delta v' + \Delta v'',$$

where

$$\Delta v' = \left(\frac{\partial v}{\partial p}\right)_S \Delta p, \quad \Delta v'' = \left(\frac{\partial v}{\partial S}\right)_p \Delta S = \frac{T}{c_p}\left(\frac{\partial v}{\partial T}\right)_p \Delta S. \tag{2}$$

Shifted components of the Rayleigh line are associated with the fluctuations $\Delta v'$, unshifted components with $\Delta v''$. The fluctuations Δp are propagated as sound waves and scatter light only at sufficiently high sonic frequencies ω which satisfy the Mandel'shtam-Brillouin condition [2]. It is evident that in a cooled liquid only the subsonic waves for which $\omega \sim |q|/T_g$ can be frozen. The $\Delta v'$ fluctuations are not frozen.

The entropy fluctuations ΔS can be represented by

$$\Delta S = \Delta S_1 + \Delta S_2 \tag{3}$$

and, correspondingly,

$$\Delta v' = \frac{T}{c_p}\left(\frac{\partial v}{\partial T}\right)_p (\Delta S_1 + \Delta S_2) = \Delta v_1 + \Delta v_2. \tag{4}$$

Here ΔS_1 are fluctuations associated with translational and rotational degrees of freedom and ΔS_2 with vibrational degrees of freedom. The resorption time of ΔS_1 is τ, the structural relaxation time of the liquid, which increases sharply at lower temperatures. For ΔS_2 the resorption time τ' is the relaxation time of the temperature conductivity [13]

$$\tau' \sim l^2/\chi, \tag{5}$$

where χ is the coefficient of temperature conductivity and l is the linear dimension of the fluctuations. Since τ' is practically independent of temperature, Δv_2 like $\Delta v'$ does not freeze. In glass we have the equilibrium values

$$\overline{(\Delta v')^2} = -kT\left[\left(\frac{\partial v}{\partial p}\right)_S\right]_T; \tag{6}$$

$$\overline{(\Delta v_2)^2} = \frac{kT^2(c_{p_2})_T}{(c_p^2)_T}\left[\left(\frac{\partial v}{\partial T}\right)_p^2\right]_T; \tag{7}$$

here c_{p_2} is the specific heat of the vibrational degrees of freedom.

The Δv_1 fluctuations associated with the structural entropy are vitrified because of the strong dependence of τ on T, at some temperature T_g. We thus obtain

$$\overline{(\Delta v_1)^2} = \frac{kT_g^2(c_{p_1})_{T_g}}{(c_p^2)_{T_g}}\left[\left(\frac{\partial v}{\partial T}\right)_p^2\right]_{T_g}; \tag{8}$$

where c_{p_1} is the structural specific heat. The intensity ratio of the central and shifted components is

$$\frac{J_c}{J_s} = \frac{\overline{(\Delta v_1)^2} + \overline{(\Delta v_2)^2}}{\overline{(\Delta v')^2}} =$$

$$= \left(\frac{c_p - c_v}{c_v}\right)_{T_g} \frac{(\partial v/\partial p)_{T_g}(c_v/c_p)_{T_g} T_g (c_{p_1})_{T_g}}{(\partial v/\partial p)_T (c_v/c_p)_T T (c_p)_{T_g}} + \left(\frac{c_p - c_v}{c_v}\right)_T \frac{(c_{p_2})_T}{(c_p)_T}. \tag{9}$$

When $T = T_g$ this formula goes over into the Landau-Placzek formula, since $c_p = c_{p_1} + c_{p_2}$.

For T_g considerable higher than T the intensity of the central component must be much greater than that of the shifted lines; this is in agreement with experiment. The central line must be narrow because the rate of resorption of ΔS_1 is very small. It is clear that V. L. Ginzburg's theory [14] of the widths of the Rayleigh line components must be amended by taking into account the resorption of ΔS_1.

We now obtain the vitrification temperature T_g of the Δv_1 fluctuations. We have the kinetic equation [15]

$$\frac{d\Delta v_1}{dt} = -\frac{1}{\tau}\Delta v_1 + A(t); \tag{10}$$

here A(t) is a purely statistical quantity.

The solution of this equation with the temperature changing at the rate $q = -|q|$ is

$$\Delta v_t = \Delta v_1^0 \exp\left[-\int_0^t \frac{dt'}{\tau}\right] + \exp\left[-\int_0^t \frac{dt''}{\tau}\right]\int_0^t \exp\left[\int_0^{t'} \frac{dt''}{\tau}\right] A(t')\,dt'; \tag{11}$$

where Δv_1^0 is the initial value of Δv_1. We shall use Chandrasekhar's method [16] for the stationary case (τ = = const). When $t \to \infty$ the fluctuation distribution approaches the Gaussian form

$$\lim_{t\to\infty} W(\Delta v_1) = \frac{1}{(T/c_p)(\partial v/\partial T)_p (2\pi k c_{p1})^{1/2}} \exp\left[-\frac{c_p^2 (\Delta v_1)^2}{2kT^2 c_{p1}(\partial v/\partial T)_p^2}\right]. \tag{12}$$

If (12) is to hold true it is necessary, when the integral on the right-hand side of (11) is divided into small parts:

$$\exp\left[-\int_0^t \frac{dt'}{\tau}\right]\int_0^t \exp\left[\int_0^{t'} \frac{dt''}{\tau}\right] A(t')\,dt' \cong \sum_{i=0}^{n} \exp\left[\int_{i\Delta t}^{t} \frac{dt'}{\tau}\right] \int_{i\Delta t}^{(i+1)\Delta t} A(t')\,dt' \tag{13}$$

that the probability of observing the different values

$$B(\Delta t) = \int_t^{t+\Delta t} A(t')\,dt' \tag{14}$$

should also obey a Gaussian distribution

$$w(B) = \frac{1}{(4\pi b \Delta t)^{1/2}} \exp(-B^2/4b\Delta t), \tag{15}$$

where

$$b = \frac{kT^2 (\partial v/\partial T)_p^2 c_{p1}}{c_p^2 \tau(t)}. \tag{16}$$

With cooling $\overline{B^2} = 2b\Delta t$, and thus $\overline{A^2}$ decreases rapidly.

Equation (13) is not rigorous in the nonstationary case which we have discussed, because in the vitrification interval the factor $\exp\left(-\int_{t'}^{t} \frac{dt''}{\tau}\right)$ can change more rapidly than A(t'). Our theory is valid only as a zero order approximation and becomes more accurate as the vitrification interval is narrowed. The characteristics of the process in this interval are not taken into account by this theory.

When (15) holds true, we have for

$$R(t) = \int_0^t \exp\left[-\int_\xi^t \frac{dt'}{\tau}\right] A(\xi)\,d\xi \equiv \int_0^t \psi(\xi) A(\xi)\,d\xi \tag{17}$$

the distribution

$$W(R) = (\pi D(t))^{-1/2} \exp(-R^2/D(t)), \tag{18}$$

where

$$D(t) = 4k \int_0^t \frac{c_{p1}}{c_p^3} \left(\frac{\partial v}{\partial T}\right)_p^2 T^2 \frac{1}{\tau(\xi)} \psi^2(\xi)\, d\xi. \tag{19}$$

We divide the integral D(t) into parts from t = 0 (beginning of cooling T = T_0) to t = t_1 (termination of cooling, T = T) and from t_1 to t (constant temperature T). In the second integral $\tau = \tau(T) = \text{const}$. The first integral can be represented by

$$4k \int_0^t \frac{c_{p1}}{c_p^3} \left(\frac{\partial v}{\partial T}\right)_p^2 T^2 \frac{1}{\tau(\xi)} \psi^2(\xi)\, d\xi =$$

$$= 4k \exp\left[-2\frac{t-t_1}{\tau(T)}\right] \int_T^{T'} \frac{c_{p1}}{c_p^3} \left(\frac{\partial v}{\partial T}\right)_p^2 T'^2 \frac{1}{|q|\tau(T')} \exp\left[-2\int_T^{T'} \frac{dT''}{|q|\tau}\right] dT'. \tag{20}$$

The quantities c_{p1}, c_p, $(\partial v / \partial T)_p^2$, T^2 are much less dependent on temperature than τ and ψ^2. The function

$$f(T') = \frac{1}{|q|\tau(T')} \exp\left[-2\int_T^{T'} \frac{dT''}{|q|\tau}\right] \tag{21}$$

has a maximum at T' = T_g, for which

$$(d\tau / dT')_{T'=T_g} = -2/|q|. \tag{22}$$

For this temperature we place $\frac{c_{p1}}{c_p^3}\left(\frac{\partial v}{\partial T}\right)_p^2 T^2$ outside of the integral, thus obtaining

$$D(t) \cong 4k\left[\left\{\frac{c_{p1}}{c_p^3}\left(\frac{\partial v}{\partial T}\right)_p^2\right\}_{T_g} \frac{T_g^2}{2} \exp\left[-\frac{2}{\tau(T)}(t-t_1)\right](1-\exp[-2\Delta\varphi(T_1,T_0)]) + \right.$$

$$\left. + \left\{\frac{c_{p1}}{c_p^3}\left(\frac{\partial v}{\partial T}\right)_p^2\right\}_T \frac{T^2}{2}\left(1-\exp\left[-\frac{2}{\tau(T)}(t-t_1)\right]\right)\right], \tag{23}$$

where

$$\Delta\varphi(T_1, T_0) = \int_T^{T_g} \frac{dT''}{|q|\tau(T'')} \to \infty, \tag{24}$$

because τ is very small for $T_0 > T_g$ and $\tau \to \infty$ for $T < T_g$ (see [12]). The final distribution (for $t - t_1 \ll \tau(T) \to \infty$) is

$$W(\Delta v_1) = (\pi D)^{-1/2} \exp[-(\Delta v_1)^2 / D], \tag{25}$$

where

$$D = \frac{2kT_g^2 (c_{p1})_{T_g}}{(c_p^3)_{T_g}} \left[\left(\frac{\partial v}{\partial T}\right)_p^2\right]_{T_g}, \tag{26}$$

. and we obtain (8).

Since according to (23) $\overline{(\Delta v_1)^2}$ has the relaxation time $\tau/2$, for $\overline{(\Delta v_1)^2}$ the vitrification condition (22) agrees with (1).

We thus see that increased intensity of light scattering in glasses results from the freezing of the fluctuations of structural entropy. When the liquid consists of anisotropic molecules we also have freezing of the orientation fluctuations, which thus contribute to the central component intensity for glasses. In inhomogeneous liquids the concentration fluctuations freeze, as the rate of their resorption through diffusion depends strongly on temperature. It is clear that this theory can be applied to viscous liquids. The Δv_1 fluctuations can be partially vitrified sooner than the entire liquid.

I am indebted to V. L. Ginzburg and V. G. Levich for valuable discussions.

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NEUTRON YIELDS FROM Rn + B, Rn + C, Rn + CaF_2, Rn + Mg Rn + Al, Rn + Si, Rn + SiO_2, Rn + GRANITE

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(Presented by Academician A. P. Vinogradov May 16, 1957)

It is known that when light elements are irradiated with alpha particles from naturally radioactive substances neutrons are produced. A fundamental distinguishing characteristic of such a neutron source is its yield, that is, the number of neutrons emitted per second unit of radioactivity. In this paper we present results of an investigation of yields and other parameters of radon sources of neutrons.

All of the neutron sources were cylindrical glass ampules 20 mm in diameter and 40 mm long. These dimensions were selected after an analysis of the yield from sources of different sizes. The ampules were completely filled with powders of the target materials, which were beryllium, boron, carbon, calcium fluoride, magnesium, aluminum, silicon, silicon dioxide and granite (porphyric granite 1.3 billion years old from the Kirovograd region).

The neutron yield was also measured without filters in ampules of borax and nonborax glass. Impurities in the pure elements and compounds did not constitute more than 1%, with the exception of amorphous boron, in which case the purity was 96%. The sources were filled with radon from the usual emanation apparatus. The radon content was determined from the external gamma radiation compared with a radium standard. Corrections were made for gamma absorption in the source and standard.

The total neutron count from the sources was obtained in two ways: by a wide-range boron counter and by the density distribution of slow neutrons in a tank of water. The boron counter with its high sensitivity and low background enabled us to obtain rapid measurements of weak neutron intensities in the presence of large gamma backgrounds. Since immediately after preparation of a radon source there is a short period during which RaA is almost in equilibrium with the radon and RaC' has still not accumulated in an appreciable quantity, the measurements obtained with the counter during this period provide a ratio between the yield of a Rn + RaA source (5.5 and 6 Mev alpha particles) and the yield of a Rn + RaA + RaC' source (7.7 Mev).

TABLE 1

Target material	Ratio to yield from Rn + Be source, in %	B in neut./sec · mC	Ratio of yield from Rn + RaA to yield from Rn + RaA + RaC in %	L, in cm	E_n, in Mev
Beryllium	100	14800±200	45,5±2.5	9.6	4.74
Boron	21.7	3220±36	54,5±2.7	8.1	3.66
Carbon	0.25	37.0±2.0	33±6.6	7.1	2.93
CaF_2	8.31	1230±15	33±2.5	6.1	2.20
Magnesium	2.85	422±6	34±3	8.1	3.66
Aluminum	2.53	374±5	20±2	5.8	1.98
Silicon	0.37	55±2.0	24±5	6.7	2.56
SiO_2	0.22	32.2±3.6	—	—	—
Granite	0.54	80.7±4.0	28±4.5	—	—
Glass					
Borax	0.92	136±8	—	—	—
Nonborax	0.74	109±7	—	—	—

Despite the large error which resulted from the limited time of measurement it was apparent that various elements differ with respect to the increase of neutron yield from RaC' alpha particles.

The second method of counting neutrons made it possible to check the range of the boron counter, to increase the total accuracy of the measurements and to estimate the average neutron energy, because the latter is linearly related to the relaxation length L, which is the reciprocal of the exponent of decline of the neutron density distribution in water [1].

As the neutron background we took the yield from ampules containing radon and filled with zinc, selenium and cadmium powders. We obtained a background of 4 ± 1.5 neutrons/sec for 1 millicurie, which is apparently due in large part to neutron emission by the short-lived radon deposit [2]; the calculated upper limit of the background which results from alpha bombardment of the ampule glass is 1-1.5 neutrons/sec per millicurie of radon.

For all sources we obtained the relative yields (as a percentage of the yield from a Rn + Be source) and the absolute yields. Only the errors of the relative measurements are included. The systematic error associated with calibration of the neutron standard does not exceed 4%.

As a check of the reproducibility of the entire process of preparing and measuring the sources we repeated the experiments with beryllium, boron, granite and glass, and obtained results which agreed within the limits of measurement errors. This serves as evidence that the yield was free of extraneous effects due to self-absorption of the alpha particles, nonuniform distribution of the radioactive material, etc.

The experimental results are given in Table 1.

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ON THE DIFFUSION OF CHARGED PARTICLES IN A UNIFORM ELECTROMAGNETIC FIELD

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In recent papers [1-3] we have examined the diffusion of molecules, neutrons and cascade particles. We have written down equations and given approximate expressions for the probability $\nu_{ij}(s, \mathbf{q}, \mathbf{u}, t, \mathbf{r}, \mathbf{v})\,d\mathbf{r}\,d\mathbf{v}$ that a particle of type i which at the instant s has the location $\mathbf{q}$ and the velocity $\mathbf{u}$ will at the time t produce a particle of type j with radius vector between $\mathbf{r}$ and $\mathbf{r} + d\mathbf{r}$ and velocity between $\mathbf{v}$ and $\mathbf{v} + d\mathbf{v}$. Using the notation and results of [3] we shall write the equation for the function V and shall give a method of solving this equation approximately when the particles whose diffusion we are investigating are charged and are moving in a uniform constant medium within a uniform electromagnetic field. We shall assume that the particles are identical and that collisions with the molecules of the medium can result not only in scattering but also in absorption and the creation of new particles; and we shall also assume that the motion of each particle between collisions is subject only to the given external field and not to the fields of the other particles. This occurs, for example, in a weakly ionized plasma which is in a stationary state, where ions of a give type collide with each other and with other ions and gas molecules, with resultant scattering, recombination and creation of new ions through impact ionization [4].

V satisfies an equation of the type of Eq. (17) of [3]. In the present paper this equation can be somewhat simplified, because the problem is more specific, as follows:

1. Since in this case the cascade consists of one kind of particle we shall have n = 1 and the subscripts i and j drop out.

2. If it is assumed without limiting generality that the magnetic field $\mathbf{H}$ is parallel to the x axis and that the electric field $\mathbf{E}$ is in the xy plane the function $\mathbf{F}$ in Eq. (1) of [3] becomes $\mathbf{E}(e/m) + \mathbf{v}(\mu e/m)\mathbf{H}$, and the integration of this equation gives for $\vec{\varphi}^s$ and $\vec{\psi}^s$ (Eq. (2) of [3])

$$\vec{\varphi}^s(t, \mathbf{r}, \mathbf{v}) = \mathbf{r} + \mathbf{h}_{s-t}(\mathbf{v}),\quad \vec{\psi}_s(t, \mathbf{r}, \mathbf{v}) = \mathbf{k}_{s-t}(\mathbf{v}),$$

where [4]

$$\mathbf{h}_t(\mathbf{v}) = \mathbf{m}_t + \mathbf{v}\cdot\mu_t,\quad \mathbf{k}_t^* \mathbf{v} = \mathbf{n}_t + \mathbf{v}\cdot\nu_t,$$

$$m_0 = \frac{1}{2}\chi_0 t^2,\quad m_1 = \frac{\chi_1}{\omega^2}(1 - \cos\omega t),\quad m_2 = -\frac{\chi_1}{\omega}t + \frac{\chi_1}{\omega^2}\sin\omega t,$$

$$n_0 = \chi_0 t,\quad n_1 = \frac{\chi_1}{\omega}\sin\omega t,\quad n_2 = -\frac{\chi_1}{\omega}(1 - \cos\omega t),$$

$$\mu_0 = t,\quad \mu_1 = \frac{1}{\omega}\sin\omega t,\quad \mu_2 = \frac{1}{\omega}(1 - \cos\omega t),$$

$$\nu_0 = 1,\quad \nu_1 = \cos\omega t,\quad \nu_2 = -\sin\omega t,$$

$$\chi_0 = \frac{e}{m}E_0,\quad \chi_1 = \frac{e}{m}E_1,\quad \omega = \frac{\mu e}{m}H\ *\ [4].$$

3. From the homogeneity of the medium we obtain

$$V(s,\ \mathbf{q},\ \mathbf{u},\ t,\ \mathbf{r},\ \mathbf{v}) = V(t-s,\ \mathbf{r}-\mathbf{q},\ \mathbf{u},\ \mathbf{v}), \quad P(t,\ \mathbf{r},\ \mathbf{v},\ \mathbf{w}) = P(\mathbf{v},\ \mathbf{w}),$$

$$p(t,\ \mathbf{r},\ \mathbf{v}) = p(\mathbf{v}),\ M(t,\ \mathbf{r},\ \mathbf{v},\ s) = \exp\left[-\left|\int_t^s p\left(\left|\mathbf{k}_\tau(\mathbf{v})\right| d\tau\right|\right]\right. = M(s-t,\ \mathbf{v}).$$

For V we obtain the integral equation

$$V(t,\ \mathbf{r},\ \mathbf{u},\ \mathbf{v}) = \delta(\mathbf{r}-\mathbf{h}_{-t}(\mathbf{v}))\,\delta(\mathbf{u}-\mathbf{k}_{-t}(\mathbf{v}))\,M(-t,\ \mathbf{v}) +$$
$$+\int_0^t\int V(t-\tau,\ \mathbf{r}-\mathbf{h}_{-\tau}(\mathbf{v}),\ \mathbf{u},\ \mathbf{w})\,P(\mathbf{w},\ \mathbf{k}_{-\tau})\,M(-\tau,\ \mathbf{v})\,d\tau\,d\mathbf{v}. \tag{1}$$

An approximate solution of this equation follows [3] in general. We shall therefore only indicate a certain difference and shall present the results immediately. The solution is obtained by integrating Eq. (23) of [3]. However, in order to obtain the initial approximation V_0 (Eq. (24) of [3]) we must generalize not the solution of the Fokker-Planck equation (Eq. (250) of [5]) for free Brownian movement as given by Chandrasekhar's Eq. (178), which was our procedure in [1-3], but an analogous solution of the Fokker-Planck equation for Brownian motion in an electromagnetic field ([5], p. 249):

$$\frac{\partial V}{\partial t} + \mathbf{v}\cdot\frac{\partial V}{\partial \mathbf{r}} + \frac{\partial}{\partial \mathbf{v}}\cdot\left[\left(\frac{e}{m}\mathbf{E} + \frac{\mu e}{m}\mathbf{v}\times\mathbf{H} - \beta\mathbf{v}\right)V\right] - \frac{\beta c^2}{2}\frac{\partial}{\partial \mathbf{v}}\cdot\frac{\partial W}{\partial \mathbf{v}} = 0. \tag{2}$$

We shall first consider the case of no electric field $\mathbf{E}$; we let

$$V_0(t,\ \mathbf{r},\ \mathbf{u},\ \mathbf{v}) = \frac{\beta^3}{\pi^3}\,g\,(1 + g'e^{-\lambda' t} + g''e^{-\lambda'' t} + \ldots)\sqrt{a^{*3}a^{3}\alpha^{*}\alpha_0\alpha_1^{2}\alpha_2^{2}}\times$$
$$\times\exp\{-\lambda t - a^{*}a\,(\mathbf{v}-\mathbf{u}\cdot\beta)^2 - (\beta\mathbf{r} - \mathbf{u}\cdot\gamma - \mathbf{v}\cdot\delta)\cdot\alpha^{*}\cdot(\beta\mathbf{r} - \tilde{\gamma}\cdot\mathbf{u} - \tilde{\delta}\cdot\mathbf{v})\}. \tag{3}$$

Here the scalar $\underline{a}$ and the tensors $\alpha, \beta, \gamma, \delta$ are functions of $\underline{t}$ which are given by

$$\frac{1}{a} = 1 - e^{-2\beta t}, \quad \beta_0 = e^{-\beta t}, \quad \beta_1 = \cos\omega t e^{-\beta t}, \quad \beta_2 = \sin\omega t e^{-\beta t},$$

$$\delta_0 = \frac{1-e^{-\beta t}}{1+e^{-\beta t}},\ \delta_1 = \frac{\beta^2}{\beta^2+\omega^2}\left(\coth\beta t - \frac{\cos\omega t}{\sinh\beta t}\right),\quad \delta_2 = \frac{\beta^2}{\beta^2+\omega^2}\left(\frac{\sin\omega t}{\sinh\beta t} - \frac{\omega}{\beta}\right),$$

$$\frac{1}{\alpha_0} = \beta t - 2\,\frac{1-e^{-\beta t}}{1+e^{-\beta t}},\quad \frac{1}{\alpha_1} = \int_0^t (\delta_1^2 + \delta_2^2)\,\beta dt,\quad \alpha_2 = 0, \tag{4}$$

$$\gamma_0 = \frac{1-e^{-\beta t}}{1+e^{-\beta t}},\quad \gamma_1 = 2\int_0^t \frac{\beta_1\delta_1 - \beta_2\delta_2}{1-e^{-2\beta t}}\beta dt,\ \gamma_2 = 2\int_0^t \frac{\beta_1\delta_2 + \beta_2\delta_1}{1-e^{-2\beta t}}\beta dt,$$

while the quantities

$$a^{*},\ \alpha^{*},\ \beta,\ \lambda,\ \lambda',\ \lambda''_1, \ldots,\ g,\ g',\ g'', \ldots \tag{5}$$

• The components of the vector $\mathbf{l}$ are denoted by $l_0 = l_x$, $l_1 = l_y$ and $l_2 = l_z$. For all of our tensors ϵ we have the relations $\epsilon_{xy} = \epsilon_{xz} = \epsilon_{yx} = \epsilon_{zx} = 0$, $\epsilon_{yy} = \epsilon_{zz}$, $\epsilon_{zy} = -\epsilon_{yz}$. We therefore characterize each tensor ϵ by the quantities $\epsilon_0 = \epsilon_{xx}$, $\epsilon_1 = \epsilon_{yy}$, $\epsilon_2 = \epsilon_{zy}$. We use $\tilde{\epsilon}$ to denote the transposed form of ϵ.

are functions of $\mathbf{u}$ but not of $\underline{t}$ and must be such that V_0 in Eq. (3) is the best possible approximation of V for large $\underline{t}$. The tensor α^* must be symmetric so that $\alpha^*_{ik} = 0$. Any further generalization which would introduce factors independent of $\underline{t}$ before β, γ, δ is impossible because this would violate the initial conditions for $t = 0$.

The solution of (2) is obtained from (3) with

$$a^* = \frac{1}{c^2}, \quad \alpha^* = \frac{1}{2c^2}, \quad \lambda = 0, \; g = 1, \; g' = g'' = \ldots = 0. \tag{6}$$

In order to obtain an equation corresponding to Eqs. (47) and (53) of [3] we here denote by $W^m(\mathbf{u}, \mathbf{v})\,d\mathbf{v}$ the probability that among the particles of the $\underline{m}$th generation in the cascade generated by a single particle with initial velocity $\mathbf{u}$ there is one particle with initial velocity between $\mathbf{v}$ and $\mathbf{v} + d\mathbf{v}$. Remembering that for $\mathbf{E} = 0$ we have $\mathbf{m}_t = 0$, $\mathbf{n}_t = 0$, $M(t, v) = e^{-p(v)|t|}$ we obtain

$$W' = \delta(\mathbf{v} - \mathbf{u}), \quad W^{m+1}(\mathbf{u}, \mathbf{v}) = \int W^m(\mathbf{u}, \mathbf{w})\, Q(\mathbf{w}, \mathbf{v})\, d\mathbf{w},$$

where

$$Q(\mathbf{u}, \mathbf{v})\,d\mathbf{v} = \frac{\int_0^\infty e^{-p(u)t}\, P(\mathbf{k}_t(\mathbf{u}), \mathbf{v})\, dt}{\iint_0^\infty e^{-p(u)t}\, P(\mathbf{k}_t(\mathbf{u}), \mathbf{v})\, dt\, d\mathbf{v}}\, d\mathbf{v}$$

is the probability of finding a particle with initial velocity between $\mathbf{v}$ and $\mathbf{v} + d\mathbf{v}$ among the particles resulting from the collision of a particle with initial velocity $\mathbf{u}$.

For the time-averaged distance traversed between two collisions by a particle with initial velocity $\mathbf{u}$, instead of $\mathbf{u}/p(u)$ [3] we obtain

$$\bar{\mathbf{r}}(\mathbf{u}) = \int_0^\infty \mathbf{h}_t(\mathbf{u})\, e^{-p(u)t}\, p(u)\, dt.$$

We denote by

$$U(u) = \int \frac{a^{*3/2}}{\pi^{3/2}} \exp(-a^* v^2)\, \frac{P(\mathbf{v}, \mathbf{u})}{\int P(\mathbf{v}, \mathbf{u})\, d\mathbf{u}}$$

the velocity distribution function of the particles for large $\underline{t}$.

Then for $\bar{\mathbf{r}}$, $\bar{\mathbf{r}}^m$, $\overline{\mathbf{r}'\mathbf{r}'}$, $\overline{\mathbf{r}'\mathbf{r}^m + \mathbf{r}^m\mathbf{r}'}$ * instead of Eqs. (50), (45), (51) and (52) of [3] we obtain

$$\bar{t} = \int \frac{U(u)}{p(u)}\, d\mathbf{u}, \quad \bar{\mathbf{r}}^m(\mathbf{u}) = \int W^m(\mathbf{u}, \mathbf{v})\, \bar{\mathbf{r}}(\mathbf{v})\, d\mathbf{v},$$

$$\overline{\mathbf{r}'\mathbf{r}'} = 2\int U(u)\, \bar{\mathbf{r}}(\mathbf{u})\, \bar{\mathbf{r}}(\mathbf{u})\, d\mathbf{u},$$

$$\overline{\mathbf{r}'\mathbf{r}^m} + \overline{\mathbf{r}^m\mathbf{r}'} = \iint U(u)\, W^m(\mathbf{u}, \mathbf{v})\, [\mathbf{r}(\mathbf{u})\, \mathbf{r}(\mathbf{v}) + \mathbf{r}(\mathbf{v})\, \mathbf{r}(\mathbf{u})]\, d\mathbf{u}\, d\mathbf{v}.$$

Using these formulas instead of Eqs. (47) and (53) of [3] as the equations which express the equality of the average displacement and average dispersion of the particles as calculated from (3) by Yang's method [6], we obtain

* The product of two tensors not separated by a sign denotes their tensor product.

$$\frac{1}{\beta}\int U(u)\,|\,\mathbf{u}\cdot\nu_{t=\infty}\,|\,d\mathbf{u}=\int U(u)\,|\,\mathbf{r}'(\mathbf{u})+\mathbf{r}''(\mathbf{u})+\ldots\,|\,d\mathbf{u}; \tag{7}$$

$$\frac{1}{4\beta^2\alpha^*\cdot(t\alpha)_{t-\infty}}=\frac{1}{6t}(\overline{\mathbf{r}'\mathbf{r}'}+\overline{\mathbf{r}'\mathbf{r}''}+\overline{\mathbf{r}''\mathbf{r}'}+\overline{\mathbf{r}'\mathbf{r}'''}+\overline{\mathbf{r}'''\mathbf{r}'}+\ldots). \tag{8}$$

For the equations which express the condition that V_0 is the best possible solution of (1) for large t, instead of Eqs. (36) and (37) of [3] we have here

$$1-\int\frac{e(v)}{p(v)-\lambda}\,d\mathbf{v}=0; \tag{9}$$

$$\int\left[\frac{3}{2}\frac{e(v)}{(p(v)-\lambda)^2}+\frac{\mathbf{v}\cdot\alpha^*\cdot\delta_{t=\infty}\cdot\widetilde{\delta}'_{t=\infty}\cdot(e(v)-\varepsilon(\mathbf{v}))\cdot\mathbf{v}}{p(v)-\lambda}\right]d\mathbf{v}=0, \tag{10}$$

where $\mathbf{e}(v)$ and $\boldsymbol{\epsilon}(\mathbf{v})$ are given by

$$e(v)=\int U(w)\,P(\mathbf{w},\mathbf{v})\,d\mathbf{w},\quad v_0^2\varepsilon_0(\mathbf{v})=\int w_0^2\,U(\mathbf{w})\,P(\mathbf{w},\mathbf{v})\,d\mathbf{w},$$
$$(v_1^2+v_2^2)\,\varepsilon_1(\mathbf{v})=\int(w_1^2+w_2^2)\,U(w)\,P(\mathbf{w},\mathbf{v})\,d\mathbf{w},\quad \varepsilon_2(\mathbf{v})=0.$$

Equations (7), (8), (9) and (10) are sufficient for the determination of a^*, α^*, β, λ. It can be seen that, just as in [3], the expressions which are obtained are independent of $\mathbf{r}$ and $\mathbf{v}$ and also of $\mathbf{u}$, thus simplifying the procedure for obtaining them.

λ', λ'', . . ., g, g', g'', . . ., are given by Eq. (56) of [3] as well as by Eq. (57) if we limit ourselves to two terms of the series within parentheses in (3). Thus all quantities in (5) have been obtained.

For a nonzero electric field $\mathbf{E}$ the solution of the Fokker-Planck equation (2) is given by (3) with the substitution

$$\mathbf{v}=\mathbf{u}\cdot\beta\to\mathbf{v}-\mathbf{u}\cdot\beta-\mathbf{b},\quad \beta\mathbf{r}-\mathbf{u}\cdot\gamma-\mathbf{v}\cdot\delta\to\beta\mathbf{r}-\mathbf{u}\cdot\gamma-\mathbf{v}\cdot\delta-\mathbf{d}, \tag{11}$$

where the vectors $\mathbf{b}$ and $\mathbf{d}$ are given by

$$b_0=\frac{\chi_0}{\beta}(1-e^{-\beta t}),\quad b_1=\frac{\chi_1}{\beta^2+\omega^2}(\omega\sin\omega te^{-\beta t}-\beta\cos\omega te^{-\beta t}+\beta),$$
$$b_2=\frac{\chi_1}{\beta^2+\omega^2}(\omega\cos\omega te^{-\beta t}+\beta\sin\omega te^{-\beta t}-\omega),$$
$$d_0=\frac{\chi_0}{\beta}\left(\beta t-2\frac{1-e^{-\beta t}}{1+e^{-\beta t}}\right),\quad d_1=\int_0^t\left[\frac{2\beta}{1-e^{-2\beta t}}(\delta_1b_1+\delta_2b_2)-(\delta_1\gamma_1-\delta_2\gamma_2)\right]dt,$$
$$d_2=\int_0^t\left[\frac{2\beta}{1-e^{-2\beta t}}(\delta_1b_2-\delta_2b_1)-(\delta_1\gamma_2+\delta_2\gamma_1)\right]dt.$$

The expressions in (4) and (6) remain the same as in the absence of an electric field. Therefore we can obtain an approximate solution V_0 of (1) in the presence of an electric field from (3) with the substitution (11) while retaining the foregoing method of obtaining the quantities in (5). When the electric field is included there are no new constants to be determined, so that the same method is applicable.

In order to have some idea of the accuracy of V_0 we can use considerations similar to those in [1].

Equation (3) and the indicated method of obtaining the quantities in (5) may be of value not only for stationary processes but also for nonstationary processes in an ionized medium, because they make it possible to calculate the transfer coefficients of the medium.

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RADIATION FROM CHARGED PARTICLES PASSING NEAR A PERFECT CONDUCTOR

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(Presented by Academician M. A. Leontovich May 18, 1957)

In connection with the development of microwave techniques and the generation of millimeter waves there has recently been increased interest in the radiation from charged particles passing near dielectrics and conductors [1-4]. The special problem of radiation from a point charge in uniform motion near a perfectly conducting sphere was examined in [5]. The general problem of radiation from charged particles passing near perfect conductors is studied here in a nonrelativistic approximation.

Let us consider the radiation from a point charge of mass m and charge e passing near a perfectly conducting surface S. We shall assume that the surface S possesses axial symmetry, so that its equations are of the form $r = h_1(s)$, $z = h_2(s)$, where s is the arc length ($-\infty < s < \infty$) with $r(s) \neq 0$, $\lim_{s \to \pm\infty} r(s) \neq 0$. The charge moves along the axis of the system from negative to positive values of s.

The exact problem is very complicated. In the present note we shall confine ourselves to a nonrelativistic approximation, in which case the problem can be divided into two parts:

I. Neglecting retardation, to integrate the equation of motion of the charge and to obtain the magnitudes of the charges and currents induced on the screen (the electric field will be regarded as electrostatic and dependent on time).

II. To compute the radiation from the system of currents obtained through solution of problem I.

Mathematically problem I reduces to the following boundary problem:

$$\Delta u = 0 \text{ in } T, \qquad u|_S = -u_0|_S, \tag{1}$$

$$m\ddot{z}_0 = -e\frac{\partial u}{\partial z}(0, z, z_0)|_{z=z_0}, \qquad \lim_{t \to -\infty} \dot{z}(t) = v_0, \tag{2}$$

where $u_0(r, z, z_0)$ is the Coulomb potential of the charge e located at the point $M_0(0, z_0)$. The function $u(r, z, z_0)$, except for the factor $4\pi e$, is the regular part g of the Green's function of Dirichlet's problem for the region T bounded by the surface S.

It will be convenient hereinafter to introduce the dimensionless coordinates

$$\rho = \frac{r}{a}, \quad \zeta = \frac{z}{a}, \quad \zeta_0 = \frac{z_0}{a}, \quad \alpha = \frac{s}{a},$$

where a is some characteristic dimension of the system such as the minimum radius of the channel or opening through which the charge e passes. The equations of S in the new coordinates are

$$\rho = \rho(\alpha) = \frac{1}{a} h_1(a\alpha), \quad \zeta = \zeta(\alpha) = \frac{1}{a} h_2(a\alpha).$$

Integrating the equation of motion (2) we obtain an expression for the velocity of the charge:

$$v(\zeta_0) = \sqrt{v_0^2 + \frac{2e^2}{am} f(\zeta_0)}; \tag{3}$$

here

$$f(\zeta_0) = \int_{-\infty}^{\zeta_0} \varphi(\varepsilon)\, d\varepsilon, \quad \varphi(\zeta_0) = -4\pi \frac{\partial g}{\partial \zeta}(0, \zeta, \zeta_0)|_{\zeta=\zeta_0}.$$

The density $\sigma(\alpha, \zeta_0)$ of the charges and $j(\alpha, \zeta_0)$ of the currents induced on S are determined by the relations

$$\sigma(\alpha, \zeta_0) = \frac{e}{a^2}\sigma_0(\alpha, \zeta_0) = \frac{e}{a^2}\frac{\partial G}{\partial \nu}(\rho(\alpha), \zeta(\alpha), \zeta_0)$$

($\nu = \frac{n}{a}$, n is the outer normal),

$$j(\alpha, \zeta_0) = \frac{ev(\zeta_0)}{a^2} j_0(\alpha, \zeta_0) = \frac{ev(\zeta_0)}{a^2}\frac{1}{\rho(\alpha)}\int_\alpha^\infty \frac{\partial \sigma_0(\varepsilon, \zeta_0)}{\partial \zeta_0}\rho(\varepsilon)\, d\varepsilon.$$

Since the law of motion of the charge and the currents which it induces on the screen have been found, it remains to compute the radiation from this system of currents. Confining ourselves to the dipole approximation, we represent the total dipole moment of the system $p(t)$ as the sum of the moments of the charge $p_1(t)$ and of the screen $p_2(t)$. For the second derivatives $\ddot{p}_1$ and $\ddot{p}_2$ we obtain

$$\ddot{p}_1 = e\ddot{z}_0, \quad \ddot{p}_2 = \frac{dJ}{dt};$$

here J is the total current in the screen

$$J = ev(\zeta_0)\chi(\zeta_0), \text{ where } \chi(\zeta_0) = 2\pi\int_{-\infty}^{\infty} j_0(\alpha, \zeta_0)\rho(\alpha)\frac{d\zeta(\alpha)}{d\alpha}d\alpha,$$

and thus

$$\frac{dJ}{dt} = \frac{e^3}{a^2 m}\varphi(\zeta_0)\chi(\zeta_0) + \frac{ev^3(\zeta_0)}{a}\psi(\zeta_0), \quad \text{где } \psi(\zeta_0) = \chi'(\zeta_0).$$

Thus

$$\ddot{p} = \frac{e^3}{a^2 m}(1 + \chi(\zeta_0))\varphi(\zeta_0) + \frac{ev^3(\zeta_0)}{a}\psi(\zeta_0),$$

and we arrive at the following formulas for the power w and the total radiation E:

$$w = \frac{2}{3c^3}|\ddot{p}|^2 = \frac{2}{3}\frac{e^2 c}{a^2}\{\gamma^2(1+\chi)^2\varphi^2 + 2\gamma(1+\chi)\varphi\psi[\beta^2 + 2\gamma f] + \psi^2[\beta^2 + 2\gamma f]^2\}, \tag{4}$$

$$E = \int_{-\infty}^{\infty} w\, dt = \frac{2}{3}\frac{e^3}{a}\left\{\gamma^2\int_{-\infty}^{\infty}\frac{(1+\chi)^2\varphi^2 d\zeta_0}{\sqrt{\beta^2 + 2\gamma f}} + 2\gamma\int_{-\infty}^{\infty}(1+\chi)\varphi\psi\sqrt{\beta^2 + 2\gamma f}\, d\zeta_0 + \right.$$

$$\left. + \int_{-\infty}^{\infty}\psi^2[\beta^2 + 2\gamma f]^{3/2} d\zeta_0\right\} = E_1 + E_2 + E_3; \tag{5}$$

here $\beta = \frac{v_0}{c}$, $\gamma = \frac{e^2}{amc^2}$. For a single electron $\gamma = \frac{r_0}{a}$, where $r_0 = \frac{e^2}{mc^2} = 2{,}8 \cdot 10^{-13}$ cm is the classical electron radius. For a "bunch" of N particles $\gamma = \gamma_N = \frac{Nr_0}{a}$.

Both dimensionless parameters β and γ in (4) and (5) are small, but by varying v_0 we shall vary the relation between them.

Let us consider the two limiting cases:

1. $\beta^2 \ll \gamma$. Then from (3) we obtain $v(\zeta_0) \cong \sqrt{2}\, c \sqrt{\gamma f(\zeta_0)}$; as a result the expression for the total radiation E becomes

$$E = \frac{4\sqrt{2}}{3} \frac{e^2}{a} \gamma^{3/2} \left\{ \frac{1}{4} \int_{-\infty}^{\infty} \frac{(1+\chi)^2 \varphi^2 d\zeta_0}{\sqrt{f}} + \int_{-\infty}^{\infty} (1+\chi) \varphi \psi \sqrt{f}\, d\zeta_0 + \int_{-\infty}^{\infty} \psi^2 f^{3/2} d\zeta_0 \right\}. \tag{6}$$

All three terms in (6) are of the same order of magnitude; therefore, it is necessary to take into account the nonuniformity of motion of the charge e with $\beta^2 \ll \gamma$ in order to calculate the total radiation. The emission spectrum in this case consists of wavelengths which satisfy $\lambda \geq \lambda_1$, where $\lambda_1 \sim \frac{a}{\sqrt{\gamma}} \gg a$.

2. $\gamma \ll \beta^2$. Then, according to (3), $v(\zeta_0) \simeq v_0 = \text{const}$, i. e., the motion of the charge can be regarded as uniform. Equation (5), when we take into account the relations

$$\frac{E_1}{E_3} \sim \left(\frac{\gamma}{\beta^2}\right)^2 \ll 1, \quad \frac{E_2}{E_3} \sim \frac{\gamma}{\beta^2} \ll 1$$

becomes

$$E \simeq E_3 = \frac{2}{3} \frac{e^2}{a} \beta^3 \int_{-\infty}^{\infty} \psi^2(\zeta_0)\, d\zeta_0. \tag{7}$$

To determine the spectrum we represent $\vec{p}(t) \simeq \vec{p}_2(t)$ by the Fourier integral

$$P(\mu) \simeq P_2(\mu) = ev_0 \Psi(\mu), \text{ where } \Psi(\mu) = \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{-i\mu\zeta_0} \psi(\zeta_0)\, d\zeta_0 \left(\mu = \frac{\omega a}{v_0}\right). \tag{8}$$

Thus in this instance we obtain principally wavelengths which satisfy $\lambda \geq \lambda_2$, where $\lambda_2 \sim \frac{a}{\beta} \gg a$.

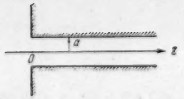

Fig. 1.

The foregoing analysis leads to the following conclusions:

1. The total radiation increases with the initial velocity as v_0^3.

2. The spectrum consists principally of wavelengths considerably greater than a, although with increasing v_0 the spectral "limit" of the radiation is shifted toward shorter waves,

3. The lower limit of applicability of the approximation for the currents (postulated constant velocity of the charge) is determined by the inequality $\frac{\gamma}{\beta^2} = 2\frac{e^2/a}{T_0} \ll 1$, where $T_0 = \frac{mv_0^2}{2}$.

The upper limit can be determined only by analyzing the relativistic problem.

Let us now consider the case in which the system is excited not by a single point charge but by a modulated beam of electrons moving with the constant velocity v_0, and assuming that their linear density on the z axis is given by

$$\eta_1(z, t) = \frac{I_0}{v_0} e^{i(\omega t - k_0 z)} \qquad \left(k_0 = \frac{\omega}{v_0}\right).$$

The assumption of constant velocity means that we neglect the interaction between the electrons and the screen and among themselves. A calculation shows that this assumption is valid if

$$\delta = 2\frac{I_0 e}{m v_0^3} \ll 1.$$

This inequality is satisfied in a very broad range of velocities and currents. Thus, for $I_0 = 0.1$ amp, $\beta = 0.1$ and we obtain $\delta = 0.01$.

The radiation of the system in this case will be monochromatic with the excitation frequency ω. In order to calculate the radiated power we must obtain the second time derivative of the dipole moment, which is given by

$$\ddot{p} \simeq \ddot{p}_2 = I_0 v_0 \Psi(\mu) e^{i\omega t},$$

where $\Psi(\mu)$ is the Fourier transform of $\psi(\zeta_0)$ (8). For the radiation resistance R and the efficiency κ of the system we obtain

$$R = 80\pi\beta^2 |\Psi(\mu)|^2 \text{ ohms;} \tag{9}$$

$$\varkappa = \frac{I_0^2 R}{I_0 V} = 0{,}98 \cdot 10^{-3} |\Psi(\mu)|^2 I_0; \tag{10}$$

here V is the potential of acceleration in volts and I_0 is the current in amperes.

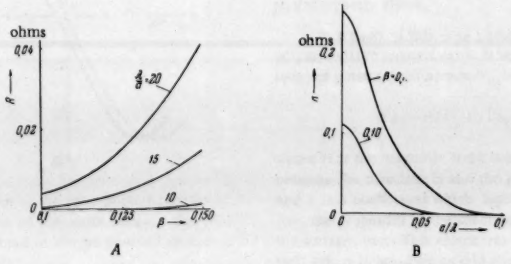

Fig. 2. A) Radiation resistance as a function of β; B) as a function of $\frac{a}{\lambda}$.

We have calculated the radiation from a beam of particles which passes from an open half space into a circular waveguide of radius a with an infinite flange (Fig. 1). $\Psi(\mu)$ is then approximately represented by

$$\Psi(\mu) \simeq 0.54\, e^{-0.63|\mu| - i\,0.09\,\mu} \sin(0.35|\mu| + 0.4).$$

Figures 2A and 2B show very strong dependence of the radiation resistance on the initial velocity, frequency and channel radius.

In conclusion the authors consider it their pleasant duty to thank R. V. Khokhlov and V. B. Braginskii for suggesting the problem and for a discussion of the results.

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ON THE ANISOTROPY OF THE EVEN OR TRANSVERSE PHOTOMAGNETIC EFFECT IN GERMANIUM SINGLE CRYSTALS

Academician I. K. Kikoin and Iu. A. Bykovskii

In [1, 2] it was shown that in addition to the ordinary photomagnetic effect in semiconductors there exists another photomagnetic effect which was called "transverse." When a specimen in the form of a flat plate is placed in a magnetic field whose direction forms some angle with the plane of the specimen, under illumination an electric field (with corresponding potential difference) is induced in the latter perpendicular to the electric field of the ordinary photomagnetic effect.

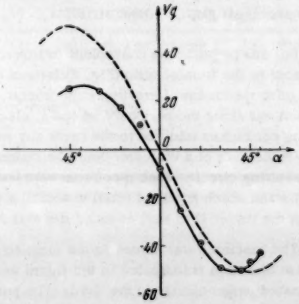

Fig. 1. Dependence of the even photomagnetic potential difference on the angle between the plane of the specimen of n-type germanium and the magnetic field.

The ordinary photomagnetic potential difference is known to be an odd function of the magnetic field, that is, its direction is reversed when the direction of the field is reversed. The potential difference of the transverse photomagnetic effect is not reversed with the field. Therefore it will be preferable hereinafter to distinguish this phenomenon by calling it the even photomagnetic effect in distinction from the ordinary odd photomagnetic effect.

In a study of this even photomagnetic effect in polycrystalline cuprous oxide it was shown [1] that the induced potential difference V_q is given by

$$V_q = AH^2 \sin 2\alpha, \qquad (1)$$

where H is the magnetic field strength, α is the angle between the direction H and the plane of the specimen and A is a coefficient which depends on the light intensity, the properties of the specimen, the condition of the surface, etc. This effect was explained as a sort of Hall effect induced by an odd photomagnetic current, thus yielding Eq. (1) immediately.

In germanium single crystals the even photomagnetic effect differs in many respects from that observed in polycrystalline cuprous oxide. Thus it was shown in [2] that in germanium single crystals the coefficient A is a complicated function of the magnetic field strength and temperature. The dependence of the even effect on the angle α was given correctly by Eq. (1). The theory of the effect, which was given in [3] and later in [4], also leads to this simple relation.

However, our recent experiments have shown that this angular dependence is not correct for single crystals. The angular dependence of the transverse photomagnetic effect which was obtained experimentally for germanium single crystals is shown in Fig. 1. The function $\sin 2\alpha$ is shown by a dashed curve. Such a marked asymmetrical angular dependence can only be explained by assuming essential anisotropy of the effect. We therefore undertook an investigation of the anisotropy of the even photomagnetic effect.

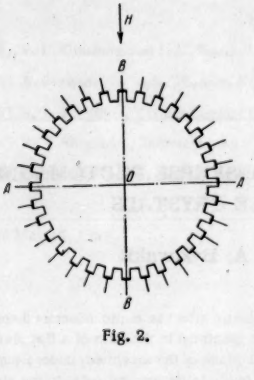

Fig. 2.

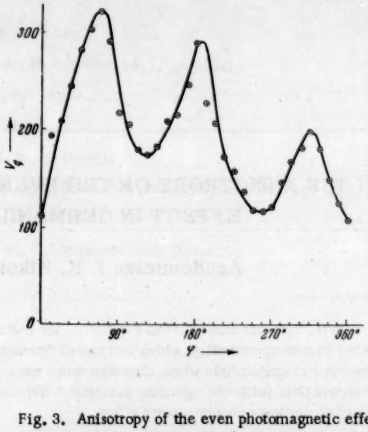

Fig. 3. Anisotropy of the even photomagnetic effect in n-type germanium. $\alpha = 45°$; H = 24,000 oersteds; the [111] axis is perpendicular to the plane of the specimen; V_q is given in arbitrary units.

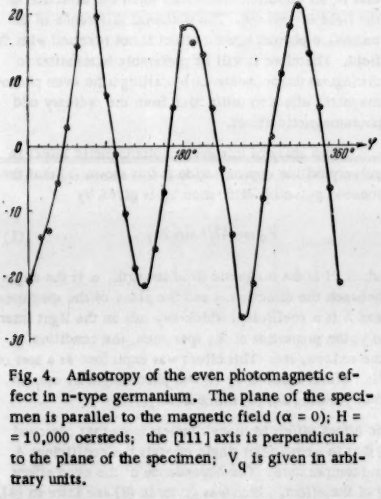

Fig. 4. Anisotropy of the even photomagnetic effect in n-type germanium. The plane of the specimen is parallel to the magnetic field ($\alpha = 0$); H = 10,000 oersteds; the [111] axis is perpendicular to the plane of the specimen; V_q is given in arbitrary units.

For this purpose it is convenient to investigate specimens in the form of disks (Fig. 2) instead of the rectangular specimens of previous experiments. The teeth cut out along the periphery served as electrodes. By using conductors soldered to the teeth any two electrodes at the ends of a diameter could be connected to the measuring circuit. Such specimens were fastened to a special stand which provided rotation around a normal through the center O as well as round the axis AA.

The specimen was placed in the magnetic field H, whose direction is indicated in the figure and was illuminated perpendicular to the field. The potential difference between two electrodes on the straight line BB was measured. The plane formed by BB and the magnetic field direction is perpendicular to the plane of the specimen. When the plane of the specimen was parallel to the field BB was also parallel to the field.

By rotating the specimen around the axis through O any two diametrically opposed electrodes can be aligned along BB. The angular separations of the electrodes are determined uniformly by the number of teeth. The number of teeth varied from 16 to 32 depending on the size of the specimen. Rotation around the AA axis changed the angle α between the plane of the specimen and the field direction (that is, between the line BB and the field direction). It is clear that for an isotropic substance the potential difference measured between any two electrodes aligned with BB for a given value of α should not depend on the angle φ of rotation around the axis through O. But experiment showed the essential dependence of the potential difference on φ. Figure 3 shows the dependence of the even photomagnetic effect on φ for a specimen of n-type germanium cut with the [111] axis perpendicular to the plane of the specimen. It is especially interesting that when the plane of the specimen was parallel to the magnetic field the even photomagnetic potential difference was not zero and also was highly anisotropic, as seen in Fig. 4.

Similar anisotropic curves were obtained with the normal to the specimen parallel to the [110] axis. In this case the curves correspond to a two-fold axis of symmetry.

The anisotropy of the even photomagnetic effect is not limited to what is stated above. It was found that the orientation of the crystal axes with respect to the magnetic field determines the dependence of the even photomagnetic potential difference on the magnetic field strength.

All of the experimental data show that the elementary theory of the even photomagnetic effect, which yields Eq. (1) and does not take the anisotropy into effect, is sharply contradicted by experiment. Thus, for example, it follows from this theory that for a given angle between the field direction and the plane of a specimen the ratio of the potential difference in the even and odd effects is proportional to the field. However, the odd photomagnetic effect is not reversed under any circumstances, whereas the even effect is reversed for certain field strengths and for certain orientations of the axes. From the same theory it follows that for $\alpha = 0$ the even effect disappears, which is also not in accord with the foregoing experimental findings.

It is thus apparent that there is no simple correlation between the even and odd photomagnetic effects and that the theory of the even effect must allow for the observed marked anisotropy. A more complete description of the experiments and the results will be published elsewhere.

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β^+-DECAY OF Ho^{160}

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O. E. Kraft and B. K. Preobrazhenskii

We have used a triple-focusing beta spectrometer [1] to investigate the holmium fraction obtained from a tantalum target that had been bombarded with 660-Mev protons. The active material was deposited in a thin layer on a 17μ cellophane film. Our measurements of the electron spectrum of holmium revealed conversion lines from transitions with the following energies: 197 ± 5; 287 ± 10; 545 ± 20; 652 ± 20; 730 ± 20; 874 ± 25; 974 ± 25 and 1315 ± 30 kev. The half-lives of these lines lie between 4.5 and 6 hours.

A study of holmium and europium fractions in our laboratory employing lens spectrometers [2] and a $\pi\sqrt{2}$ spectrometer [3], scintillation methods employed in the laboratory of the Radium Institute of the Academy of Sciences USSR [4] and the work reported in [5, 6] have shown that the conversion electrons corresponding to the following transition energies: 196, 298, 539, 648, 730, 876, 967 kev, belong to the isotope Ho^{160} with a half-life of 5.3 hours. Thus it can be assumed that the conversion spectrum of our specimen was entirely due to Ho^{160}, with the exception of the line $h\nu$ = 1315 kev, which no one has yet observed.

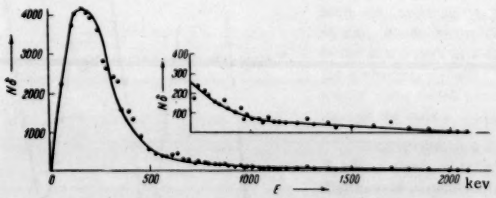

Fig. 1. Positron spectrum of Ho^{160}.

Thus far the literature contains no discussions of positrons in the radiation from Ho^{160}. Yet a study of the positron spectrum of this isotope could provide information about the mass difference between Ho^{160} and Dy^{160}. We have measured the positron spectrum of Ho^{160}. These measurements showed that the spectrum is a complex curve (Fig. 1). Assuming that this spectrum results from β^+-decay, we constructed the corresponding Curie plot shown in Fig. 2. This plot shows clearly the presence of four components in the β^+-spectrum. The energy limits of the components are 1900 ± 100; 970 ± 30; 570 ± 20 and 300 ± 50 kev and the relative intensities 0.14 : 0.26 : 1 : : 0.47. We cannot positively affirm the independent existence of the last soft component of the spectrum, because in this energy region there are large corrections for scattering by the apparatus.

To determine the half-lives of all components we observed the change of the counting rate of the positrons during 30 hours and established that the half-life of 5.6 ± 0.7 hours is the same for all components. An analysis of the results raises two questions:

1. To which holmium isotope does the positron emission belong?

2. What is the origin of these positrons – do they really result from β^+-decay or are they pair conversion positrons of the corresponding transitions?

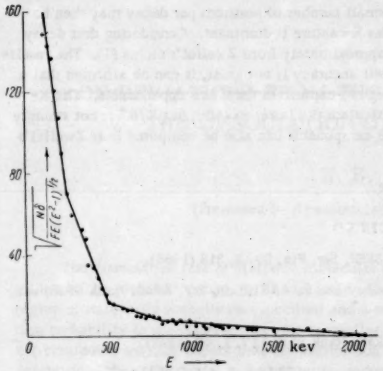

Fig. 2. Curie plot of the β^+-spectrum of Ho^{160}.

In considering the first question it must be noted that the value of the positron half-life favors the Ho^{160} hypothesis. But this does not exclude the possibility that the positron spectrum is due to some unknown Ho^X with a half-life close to 5 hours. If this unknown Ho^X is a small fraction of the amount of Ho^{160} its conversion spectrum cannot be seen against the background of the strong Ho^{160} conversion spectrum. In order to decide to what value of A the positron spectrum belongs we performed two additional experiments.

It is known from the literature that Ho^{160} is produced by electron capture in Er^{160} with a half-life of about 30 hours. The first control source was obtained by separating holmium from an erbium fraction 45 hours after its preparation. Measurements of this source gave the same results as previously – the same conversion spectrum, the same half-life and the same components of the positron spectrum. In addition, the ratio of the area of the total positron spectrum to the area of the conversion line with $h\nu$ = 730 kev, which certainly belongs to Ho^{160}, was also the same. The ratio $S_{\beta^+}/S_{e^-_{730}}$ is 1.6 ± 0.2 for both sources.

Hence it follows that:

1) since a β^+-spectrum with a half-life of 5.6 hours was observed for a source separated from Er after 45 hours, the hypothetical Ho^X like Ho^{160} is obtained from Er^X;

2) the fact that the ratio $S_{\beta^+}/S_{e^-_{730}}$ was the same for both sources shows a remarkable similarity of the half-lives not only of Ho^{160} and Ho^X but also of Er^{160} and Er^X.

In the second supplementary experiment we used an erbium fraction which lay for 110 hours after separation; the results were similar. The ratio $S_{\beta^+}/S_{e^-_{730}}$ in this case was 1.7± 0.2, which fully agrees with the previous measurements within the limits of error. If the erbium fraction contained besides Er^{160} some amount of Er^X with a half-life differing from that of Er^{160} by only 3 hours, then after 110 hours this difference would cause $S_{\beta^+}/S_{e^-_{730}}$ to change by 35%. Our experiments lead to the conclusion that the positron emission comes from holmium with mass 160.

If it is postulated that the positrons which we detected result from pair conversion the corresponding transitions must be quite intense and we should have observed strong conversion lines. In the region from 1 to 2.5 Mev we observed only one weak line at 1315 kev. The bulk of the positrons (~ 50%) have a limit at about 550 kev. Even a corresponding E1 conversion line or lines should be 100 times stronger than the observed 1315 line. Similarly it can be proved that the most energetic positrons do not result from pair production. It can therefore be assumed that most of the positrons come from the β^+-decay of Ho^{160}. Assuming that the hard part of the positron spectrum belongs to the continuous β^+-spectrum, we arrive at the conclusion that the Ho^{160}-Dy^{160} mass difference is not less than 2920 ± 100 kev.

TABLE 1

Energy limit	% decays in β^+-branch	log τf
E_1 = 300 kev	0.11	4.9
E_2 = 570 "	0.25	5.9
E_3 = 970 "	0.06	7.6
E_4 = 1900 "	0.04	9.2
Total	0.46	

For the purpose of determining the intensity of the positron spectra (the number of positrons per single decay) we can use our measured ratio $S_{\beta^+}/S_{e^-_{730}}$ = = 1.6 ± 0.2. By using the ratio e_{730}/e_{60} = 0.0029, which was measured by Grigor'ev et al. [3], we obtain β^+/e^-_{60} = = 0.46%.

Since the 60-kev transition is isomeric it is strongly converted and the number e^-_{60} practically coincides with the number of decays. (We neglect possible but unestablished beta process of Ho^{160} from a 5-hour level.)

Thus we obtain the intensities and values of $\log fT$ for the β^+-spectrum (Table 1).

Transitions to the upper levels are thus allowed. The small number of positrons per decay may then be surprising, but results from the fact that for low decay energies K-capture is dominant. Considering that decay to the two upper levels is allowed, we can determine K/β^+ approximately from Zweifel's tables [7]. The resulting values of ~ 6000 and 450 are very large and, although their accuracy is not great, it can be affirmed that a considerable fraction of all Ho^{160} transformations occurs through K-capture to these two upper levels. The K-capture intensity to lower levels is small. It is difficult to calculate the latter exactly, but K/β^+ is not strongly dependent on forbiddenness; therefore the intensities of these components can also be computed from Zweifel's tables, giving −1.7 and 0.12% of the decays.

The authors are grateful to N. Popova for her assistance.

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ON DEVIATIONS FROM THE BOLTZMANN DISTRIBUTION IN THE DISSOCIATION OF DIATOMIC MOLECULES

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(Presented by Academician V. N. Kondrat'ev April 12, 1957)

The dissociation rate of diatomic molecules in an $AB \rightarrow A + B$ reaction due to "dual" collisions is determined by the probability that the AB molecule will make a transition from a bound vibrational state into the region of states with a continuous spectrum and a distribution function of vibrational energy. When the transition probability in the discrete spectrum is smaller than the probability of dissociation from upper levels into the continuum we may expect that the distribution will deviate from the Boltzmann form. For the purpose of determining the influence of such departures from the Boltzmann distribution on the dissociation rate we make the following simplifying assumptions.

We postulate that the diatomic molecules AB which are able to dissociate are in a thermal bath consisting of molecules C and that the concentration of AB molecules is considerably smaller than that of the C molecules. Then the collisions among AB molecules can be neglected. As a result of collisions between the AB and C molecules the AB molecules begin to dissociate, but because of the small AB concentration the equilibrium distribution of the C molecules will be maintained with great accuracy.

We postulate furthermore that the probability of an inelastic collision between AB and C with a transition of AB from the quantum state n to the state n' is proportional to the square of the matrix element of the vibrational coordinate of AB for the corresponding transition. This is justified when the amplitude of the vibrations of AB is smaller than the range of intermolecular forces [1] and also when the vibrational amplitude is small compared with $(\Delta k)^{-1}$, where Δk denotes the change of the wave vector k with respect to the motion of AB and C in an inelastic collision [2].

As a model of the diatomic molecule we take a quantum system with $n_0 + 1$ discrete and approximately equidistant levels of negative energy and a continuous spectrum of positive energy. The wave functions of the discrete spectrum of the appropriate potential well are, strictly speaking, different from the wave functions of a harmonic oscillator, but the transition probabilities can be assumed to be identical to a sufficient degree of accuracy. This results from the fact that for a strongly anharmonic Morse oscillator the matrix elements for transitions between adjacent levels differ little from the corresponding harmonic oscillator matrix elements, and that the matrix elements for transitions between nonadjacent levels are considerably smaller than for transitions to adjacent levels [3]. With regard to transitions from the discrete to the continuous spectrum a calculation shows that for the Morse potential the matrix element of the coordinate is greatly reduced with increasing transition energy; it can therefore be assumed that a transition to the continuous spectrum takes place only from an upper level. Thus relaxation and dissociation of a system of oscillators in a heat reservoir is characterized by the probabilities of transitions to adjacent levels, $P_{n,n+1} = P_{01}(n + 1)$ [1], for each collision and by the probability of dissociation from the last level, P_{∞}.

If $x_n(t)$ denotes the distribution function, i. e., the probability of finding an AB molecule on the nth level, the time dependence of x_n is determined by the following set of equations in which the activation and deactivation probabilities are related through the ratio $\frac{P_{n,n+1}}{P_{n+1,n}} = \exp\left(-\frac{\hbar\omega}{kT}\right) \equiv \alpha$:

$$dx_n/dt = ZP_{10}\{\alpha n x_{n-1} - [n + (n+1)\alpha]\, x_n + (n+1)\, x_{n+1}\},$$

$$n = 0, 1, 2, \ldots, n_0 - 1,$$
$$dx_{n_*}/dt = ZP_{10}\{\alpha n_0 x_{n_*-1} - n_0 x_{n_*}\} - ZP_{\infty} x_{n_*}. \tag{1}$$

Here Z denotes the number of collisions between AB and C molecules per second with $N_C = 1$ cm^{-3}. When $n_0 \to \infty$ this system is transformed into the system of levels for a harmonic oscillator which was studied in [4].

The solution of (1) will be obtained by superposition of the terms $x_n = l_n(\mu) \exp(ZP_{10}\mu t)$. Then the coefficients μ will be the eigenvalues of the matrix of (1). We denote these eigenvalues by $\mu^{(0)}, \mu^{(1)}, \mu^{(2)}, \ldots$ in order of increasing absolute value; they determine the rates of relaxation and dissociation. The dissociation rate constant is given by

$$k = -\left(\frac{d}{dt}\sum_{n=0}^{n_*} x_n(t)\right) \Big/ \sum_{n=0}^{n_*} x_n(t) \tag{2}$$

and generally depends on time; but when $t \gg |ZP_{10}\mu^{(1)}|^{-1} = t_0$ it tends toward the constant limit

$$k = -ZP_{10}\mu^{(0)}. \tag{3}$$

The eigenvalues and the eigenvectors of the matrix $\hat{B}$ of (1) can be found by using the Lifshits theory of degenerate regular perturbations [5]. Some slight modifications are needed because the zero approximation matrix B^0 is non-Hermitian. B^0 is taken to be the matrix of (1) as $n \to \infty$; its eigenvalues and eigenvectors are given by the equations [4]

$$\mu_0^{(m)} = (\alpha - 1)m, \quad m = 0, 1, 2, \ldots;$$
$$l_k^0(\mu_0^{(m)}) = \alpha^{m/2}\sqrt{1-\alpha}\,F(-k, m+1, 1-\alpha), \tag{4}$$

where $F(a, b, c, z)$ is the hypergeometric function. The vectors $l_k^0(\mu_0)$ are orthonormalized with the weights α^{-k}, i. e.,

$$\sum_{n=0}^{\infty} l_n^0(\mu_0)\,\alpha^{-n}\,l_n^0(\mu_0') = \delta_{\mu_*\mu_0'}. \tag{5}$$

The perturbation matrix $\hat{\Lambda} = \hat{B} - \hat{B}^0$ must eliminate transitions between the n_0 and $n_0 + 1$ levels and introduce the dissociation probability P_∞ from the n_0 level. Its nonzero elements are

$$\Lambda_{n,n_*} = -P_\infty/P_{10} + \alpha(n_0+1), \quad \Lambda_{n_*,n_*+1} = -(n_0+1),$$
$$\Lambda_{n_*+1,n_*} = -\alpha(n_0+1), \quad \Lambda_{n_*+1,n_*+1} = (n_0+1).$$

For solution of the set of equations

$$\sum_k (B_{ik} - \mu\delta_{ik})\,l_k(\mu) \equiv \sum_k (B_{ik}^0 - \mu\delta_{ik} + \Lambda_{ik})\,l_k(\mu) = 0 \tag{6}$$

the matrix $(\hat{B}^0 - \mu\hat{E})$ must be diagonalized; then the order of the determinant of (6), from which the eigenvalues μ are obtained, is reduced to the order of the matrix $\hat{\Lambda}$. Following [5], we obtain

$$(\hat{B}^0 - \mu\hat{E})^{-1}_{ik} = \sum_{\mu_0} \frac{f_i^0(\mu_0) f_k^0(\mu_0)}{\mu_0 - \mu} \alpha^{-k} \equiv T_{ik}\alpha^{-k}. \tag{7}$$

The summation here is taken over all unperturbed eigenvalues μ_0. Multiplying (6) on the left by $(\hat{B} - \mu\hat{E})^{-1}$ we obtain the equations

$$\sum_k \left(\delta_{ik} + \sum_j T_{ij}\alpha^{-j}\Lambda_{jk}\right) l_k(\mu) = 0, \tag{8}$$

in which the T_{ij} depend on μ. Assuming that $\mu^{(0)} \ll \mu_0^{(1)}$, and using the expression for $F(-n, m+1, 1, 1-\alpha)$ as the sum $\alpha^n \sum_{\nu=0}^{\infty} (1-1/\alpha)^\nu \binom{n}{\nu}\binom{m}{\nu}$, we can calculate $T_{ij}(\mu^{(0)})$:

$$T_{ij}(\mu^{(0)}) = \alpha^j\left[\alpha^j\lambda^{(0)} - \frac{1}{(j+1)(1-\alpha)}\right], \quad j \leqslant i; \qquad T_{ji}(\mu^{(0)}) = T_{ij}(\mu^{(0)}), \quad i \leqslant j. \tag{9}$$

Here we have made the substitution $\lambda^{(0)} = (\alpha - 1)/\mu^{(0)}$ and assume that $i, j \gg 1$. The magnitude of $\mu^{(0)}$ is obtained by setting the determinant of (8) equal to zero, resulting in

$$\begin{vmatrix} 1 + \dfrac{P_\infty}{P_{10}(1-\alpha)(n_0+1)} - \dfrac{P_\infty}{P_{10}}\alpha^{n_0}\lambda^{(0)} & 0 \\ \alpha\left[1 + \dfrac{P_\infty}{P_{10}(1-\alpha)(n_0+1)} - \dfrac{P_\infty}{P_{10}}\alpha^{n_0}\lambda^{(0)}\right] + \dfrac{\alpha}{1-\alpha}\Omega\left(\dfrac{1}{n_0^2}\right) & \dfrac{1}{1-\alpha}\Omega\left(\dfrac{1}{n_0^2}\right) \end{vmatrix} = 0.$$

Hence, taking (3) into account, we obtain

$$k = ZP_{10}(1-\alpha)\frac{P_\infty\alpha^{n_0}}{P_{10} + P_\infty/(1-\alpha)(n_0+1)}. \tag{10}$$

Now solving (8), we obtain the distribution function $x_n(t)$ for $t \gg t_0$:

$$x_n(t) = A\alpha^n\left\{1 - \alpha^{n_0-n}\left(\frac{n_0+1}{n+1}\right)\frac{P_\infty}{P_\infty + P_{10}(1-\alpha)(n_0+1)}\right\}\exp(-kt). \tag{11}$$

The normalization constant A must be determined from the initial distribution. Physically it is the fraction of the original number N_0 of molecules which did not dissociate in the time $t \approx t_0$. An estimate shows that if the AB molecules were originally at a lower temperature the number ΔN which dissociated in $t \approx t_0$ satisfies $\frac{\Delta N}{N_0} \leqslant \frac{k}{ZP_{10}(1-\alpha)}$. If at the initial instant the molecules were on a level with the energy E, then $\frac{\Delta N}{N_0} \leqslant$ $\leqslant \exp[E - n_0\hbar\omega/kT]$.

It follows from (11) that for small $\frac{P_\infty}{P_{10}}$ the distribution function $x_n(t)$ differs little from the Boltzmann function $x_n = A\alpha^n$; for large $\frac{P_\infty}{P_{10}}$ the equilibrium distribution is disturbed so that the probability of finding an AB molecule on an upper level is close to zero. From the factor before the exponential in (11) we can obtain the energy region ϵ in which the Boltzmann distribution is essentially changed, i. e., the value of ϵ at which the second term in the brackets is reduced by the factor e. Setting $\frac{P_\infty}{P_{10}} \gg (1-\alpha)(n_0+1)$ we obtain $\epsilon = \frac{DkT}{(D-kT)}$ for $\epsilon \ll D$, where ϵ is reckoned from the boundary of the continuous spectrum and D denotes the dissociation energy which, as follows from the model, satisfies $n_0\hbar\omega \leqslant D \leqslant (n_0+1)\hbar\omega$.

The dissociation rate constant $\underline{k}$ in (1) can be written as

$$k = k_0 \frac{p \exp(-D/kT)}{1 + p \frac{\hbar\omega}{D} \Big/ \left[1 - \exp\left(-\frac{\hbar\omega}{kT}\right)\right]}, \tag{12}$$

where $k_0 = ZP_{10}(1-\alpha)$ and $p = \frac{P_{\infty}}{P_{10}}$.

The earlier assumptions $n_0 \gg 1$, $t \gg t_0$, $\mu^{(0)} \ll \mu_0^{(1)}$ now become $\frac{D}{\hbar\omega} \gg 1$, $\frac{(1-\alpha)D}{\hbar\omega \exp(-D/kT)} \ll 1$. These inequalities are practically always fulfilled. According to [4] the quantity $\frac{1}{k_0}$ is the relaxation time of the energy of harmonic oscillators in a heat bath.

When the equilibrium distribution is only slightly disturbed, that is, for $p \ll \frac{(1-\alpha)D}{\hbar\omega}$, (12) gives $k = ZP_{\infty}\exp(-D/kT)$. The dissociation rate constant is proportional to the probability of transition from the last discrete level to the continuous spectrum. Since $P_{\infty} \ll \frac{P_{10}(1-\alpha)D}{\hbar\omega} \leq 1$, the factor before the exponential is in this case much smaller than the number of collisions.

The maximum of $\underline{k}$ is attained when $p = \frac{P_{\infty}}{P_{10}} \gg \frac{(1-\alpha)D}{\hbar\omega}$, which is usually fulfilled. Then the rate constant is given by

$$k = k_0 (1-\alpha) D / \hbar\omega \exp(-D/kT) \tag{13}$$

and is not determined by the probability of a transition to the continuous spectrum but by the probability P_{01}. The magnitude of k_0 depends on the temperature, the magnitude of the transferred quantum $\hbar\omega$ and the range of the intermolecular forces. For various molecules at ordinary temperature k_0 lies between the limits $Z \cdot 10^{-2}$-$Z \cdot 10^{-10}$. Since we usually have $\alpha \ll 1$, $\frac{D}{\hbar\omega} \approx 10$-$10^2$, the maximum value of the factor before the exponential in (13) is of the order of the number of dual collisions.

The principal shortcoming of our model is the neglect of anharmonicity, i. e., constancy of $\hbar\omega$. In actuality the magnitude of the vibrational quantum in diatomic molecules is reduced with an increasing number of the vibrational level. Since k_0 increases as $\hbar\omega$ is reduced this must result in an increase of the factor before the exponential compared with that for the assumed model. The experimental values of this factor for the reaction AB + C ⟶ A + B + C often exceeds the number of collisions by two or three orders of magnitude.

I am deeply grateful to Professor N. D. Sokolov for a discussion.

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ABSORPTION SPECTRA FROM COLD LIGHT SOURCES

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The usual experimental arrangement for the study of absorption spectra is shown in Fig. 1. Here a is the light source, b is the cell containing the substance to be investigated and c is the receiver of the radiation. When all elements of the scheme are at the same temperature the energy flux which reaches the receiver from the source is fully balanced by the flux emitted by the receiver and the latter records nothing. When the source temperature is above the receiver temperature the flow of radiant energy in the positive direction (from the source to the receiver) exceeds the energy flow in the negative direction and the receiver, being heated, gives positive readings. If at the point a we place a cooled body with $T_s < T_r = T_c$ the energy reaching the receiver will be less than the energy emitted by it so that it will cool off and give negative readings. The receiver readings will depend not only on the temperatures of the light source and receiver but also on the characteristics of the cell and especially on the absorption coefficient of the investigated substance. This fact, which has been noted in [1], can easily be used to obtain absorption spectra from cold sources of light.

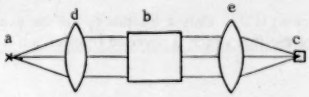

Fig. 1.

Figure 2 is the trace of a nitrobenzene spectrum from an IKS-11 infrared spectrophotometer in the usual setup. The cell, which was at room temperature, was 0.09 mm wide. The gap width was 1.9 mm. The thin line represents the zero level. Curve a represents the spectrum of a source of positive radiation (an electric furnace +44°); curve b is the spectrum of a negative source. The negative source was a brass rod cooled by liquid air to −100°. Curves c and d are the spectra of the first and second sources when the cell contained nitrobenzene. Figure 3 shows the traces for toluene obtained under similar conditions. From Figs. 2 and 3 we can see that curves c and d exhibit good symmetry with respect to the zero line; the positions of the absorption bands are identical. This result shows the possibility and sometimes the advisability of making practical use of sources of negative radiation flux. It must be remembered that the shapes of curves c and d depend not only on the absorption coefficient of the substance in the cell but also on the temperature of the furnace and refrigerator, respectively. At different temperatures the symmetry of the curves would be destroyed.

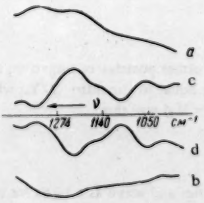

Fig. 2.

A check showed that Bouguer's law is obeyed for negative flux. The absorption coefficients calculated from the spectra for positive and negative flux are identical. These results show that in the infrared region it is absolutely necessary to take into account the thermal emission of the receiver and of the cell. In some instances the radiant energy passes from the receiver through the cell to the light source, and yet the information obtained from the instrument can be used to calculate the absorption spectrum of the investigated material.

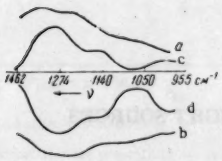

Fig. 3.

The receiver readings are generally proportional to the difference $E_i - E_r$, where E_i is the spectral intensity of the radiation incident on the receiver and E_r is the spectral intensity of the radiation emitted by the receiver. $E_i = E_1 + E_2$, where the first part represents the radiation of the light source and the second part represents the thermal emission of the substance in the cell.

Let us assume for simplicity that the light source is a black body with the emissive power U_s. Some fraction (α) of the light source energy enters the cell. The value of α is determined by the transmission of the focusing system and the dimensions of the source, cell and receiver. In the absence of the cell all of the energy would reach the receiver. In the presence of the cell the spectral intensity of the flux at the receiver is $E_1 = (1 - R)\,\alpha U_i e^{-kl}$, where R is the reflection coefficient of the investigated substance, l is the thickness of the substance involved and k is its absorption coefficient.

E_2 can easily be computed from Kirchhoff's law. The spectral intensity of the flux emitted by the cell is $(1 - R)\,A_c U_c = (1 - R)\,(1 - e^{-kl})\,U_c$, where U_c is Planck's function for $T = T_c$. The receiver is reached by only a part of this flux $E_2 = \beta(1 - R)\,(1 - e^{-kl})\,U_c$. The coefficient β is determined from the geometry of the system.

The emissive power of the receiver (a black body in our case) is U_r. Only a fraction (γ) of the energy is emitted in the direction of the cell, and the spectral intensity of the flux which is reversed in direction is $R_\gamma U_r$, and thus $E_r = (1 - R)\gamma U_r$.

The receiver readings are

$$Y = b(E_i - E_r) = b\,[\alpha\,(1 - R)e^{-kl}U_s + \beta(1 - R)\,(1 - e^{-kl})\,U_c - \gamma\,(1 - R)\,U_r]\,, \qquad (1)$$

where b is a constant of proportionality. When the system is in full thermodynamical equilibrium the light source, cell and receiver have the same temperature, and there is no flow of energy between these elements. In this case $Y = 0$. Assuming in (1) that $U_s = U_c = U_r$ and that this equality is valid for all values of kl, we obtain

$$\alpha = \beta = \gamma. \qquad (2)$$

We obtained (2) from general thermodynamic considerations. It could also be obtained from simple considerations of photometry and geometric optics. Inserting (2) in (1), we finally obtain

$$Y = \alpha b\,(1 - R)\,[U_s e^{-kl} + (1 - e^{-kl})\,U_c - U_r]\,. \qquad (3)$$

Equation (3) embodies many special cases of practical importance. For example, it must be taken into account when measuring the temperature dependence of the absorption coefficient or when measuring temperatures by spectral reversal. We are here interested only in the most frequent case $U_r = U_c$. In this instance

$$Y = \alpha b\,(1 - R)\,(U_s - U_r)\,e^{-kl}\,. \qquad (4)$$

Depending on the temperatures of the source and receiver Y can be either positive or negative, but independently of this the transmission P of the system is always positive, being equal to the ratio Y/Y_0, where Y_0 is the reading of the receiver for an empty cell. The value of Y_0 is $\alpha b\,(U_s - U_r)$ and thus

$$P = (1 - R)\,e^{-kl}. \qquad (5)$$

Equations (4) and (5) are in good agreement with experimental findings and serve as a basis for utilizing sources of negative radiation in the measurement of absorption coefficients.

All of the effects described here are appreciable only in the infrared region or at very high temperatures. In the visible region the emission of the receiver is negligible and the possible negative flow cannot be recorded.

It follows from (3) that measurement of the absorption coefficient $\underline{k}$ is also possible without a light source if the cell is heated or cooled. This method is based on measurement of the thermal emission of the substance in the cell.

Equations (3) and (4) can also be used to describe the negative opto-acoustic effect recently discovered by Veingerov, Gerlovin and Pankratov [2].

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THE EFFECT OF THIN MERCURY FILMS ON THE STRENGTH OF METAL SINGLE CRYSTALS

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The reduction of the surface energy of a metal through adsorption of surface-active substances results in decreased strength [1, 2]. The very general aspects of adsorptive strength reduction have been studied in detail for metals which had adsorbed organic surface-active materials under various conditions [2-4]. However, metallic liquids can induce considerably greater reduction of strength by lowering the solid-to-liquid interphase surface energy to almost zero. This effect explains the results obtained by S. T. Kishkin, Ia. M. Potak and several other investigators [5-7] of the reduced strength of metals at high temperatures, affected by low-melting metals.

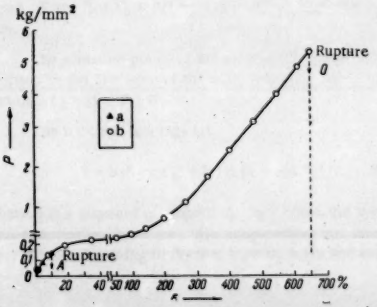

Fig. 1. Dependence of the true stress $P = P_0\left(1 + \frac{\epsilon}{100}\right)$ on the relative stretching ϵ for amalgamated (a) and unamalgamated (b) zinc single crystals. $\chi_0 = 40°$, $\lambda = 40°$. The stretching rate was 15% per minute.

Rebinder, Likhtman and Kochanova [8] have shown that adsorptive strength reduction under the influence of surface-active liquid metals saturated with the metal to be investigated is very pronounced in perfect single crystals.

These experiments were performed at high temperatures. It was thus of interest to study the effect of a surface-active metallic coating on the strength of single crystals at room temperature. We used single crystals of zinc, tin, cadmium and lead 99.99% pure, of 0.5 mm diameter and ~ 10 mm length. The surface-active substance was mercury which was deposited in a thin film by dipping the specimens in a solution of $Hg_2(NO_3)_2$ followed by washing in water. The mercury covered the crystals with a uniform layer about 0.1 μ thick and was quickly saturated with the investigated metal. The extremely small amounts of deposited mercury and the relatively low solubility of the metals in mercury (2.2% for Zn and 0.6% for Sn at 20°) prevented any appreciable corrosive effect of the mercury.

An investigation of the strength of the amalgamated single crystals when stretched at a constant rate showed that the strength of zinc and tin single crystals coated with mercury is several times less than that of unamalgamated specimens (Figs. 1 and 2).

This sharp reduction of the strength is evidently associated with considerable lowering of the surface tension at the metal-to-mercury boundary and by the reduced energy required to form a new surface crack leading to breaking. The development of such cracks in a slip plane is apparently aided by surface migration of mercury along developing lattice defects.

An investigation of axial sections of amalgamated zinc crystals after deformation showed that the cracks can form both on the surface and in the body of the crystal. This is shown by an examination of the axial slices (Fig. 3, I). The formation of body cracks can be associated with considerable diffusion of the mercury into the zinc at room temperature, followed by reduced surface tension on the internal surfaces of separation.

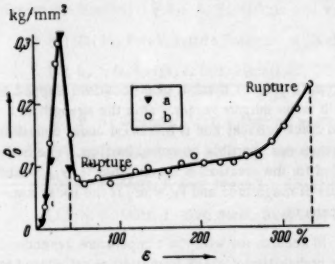

Fig. 2. Dependence of the stress P_0 on the relative stretching ϵ for amalgamated (a) and unamalgamated (b) tin single crystals. $\chi_0 = 58°$, $\lambda = 74°$. The stretching rate was 15% per minute.

A study of the temperature dependence of the adsorptive effect of mercury showed that by raising the temperature to 160°C one eliminates the described catastrophic brittleness of zinc single crystals, thus completely restoring the plasticity and strength (Fig. 3,II). In a 40° temperature interval (120-160°) of the transition from the brittle to the plastic state "partial embrittlement" is observed (Fig. 3, II) when the crystal is covered with a number of more or less deep half-opened cracks around which the material is strongly consolidated. These pronounced temperature effects are fully reversible; thus they are not associated with removal of the mercury by evaporation. A reduced deformation rate produces effects analogous to the raising of the temperature.

The mechanism of crack formation, which has been thoroughly studied in several papers [9-11], is usually associated with the clustering of dislocations

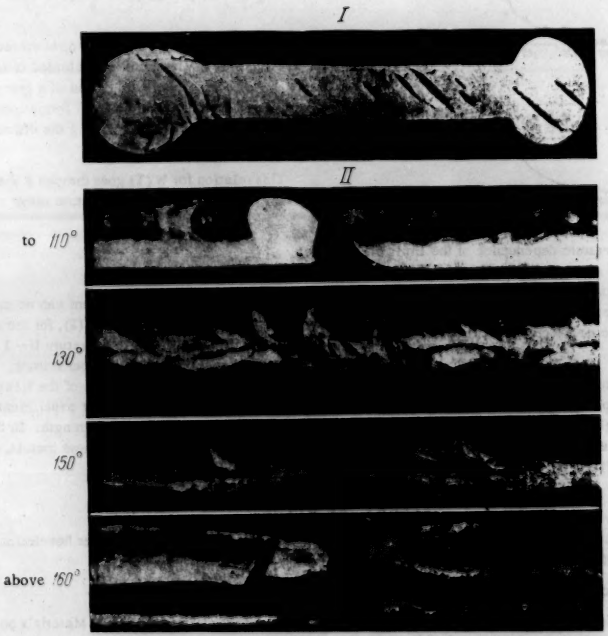

Fig. 3. I) Axial slice; II) destruction of amalgamated zinc single crystals at different temperatures. Stretching rate 15% per minute.

near various obstacles. The stress P_c which is sufficient for crack formation at the head of such a cluster was calculated by Stroh [9] to be

$$P_c \approx \frac{12\sigma_{12}}{nb}.$$

Here σ_{12} is the surface tension at the boundary of the single crystal; n is the number of dislocations stopped at the obstacle and depends on the penetrability of the obstacle; b is the Burgers vector. Thus the strength of a body is reduced with lowered surface tension on freshly-formed crack surfaces and is increased under conditions which prevent the accumulation of large avalanches of dislocations near possible potential barriers. The body is brittle when the yield point P_f is higher than the stress P_c needed for the creation of cracks, that is, $\sigma_{12} < \alpha b P_1$, where α is a theoretically estimated constant near unity in order of magnitude and $P_1 \approx nP_f$ is the local stress needed to overcome the obstacle for propagation of the displacement.

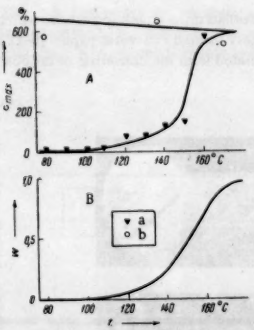

Fig. 4. A) Temperature dependence of the limit of elongation before rupture of (a) amalgamated and (b) unamalgamated zinc single crystals; B) temperature dependence of the probability that dislocations in an amalgamated zinc single crystal will overcome an obstacle.

In accordance with the temperature dependence of the probability of crack formation as calculated by Stroh [9] the probability that a cluster of dislocations will overcome an obstacle, that is, the probability of plastic flow, can be expressed by

$$W(T) = 1 - \exp\left(-\frac{\nu l t_0}{b} e^{-U/kT}\right), \qquad (1)$$

where ν is natural frequency of thermal vibrations of the atoms; l is the length of the retarded dislocation front; t_0 is the time of accumulation of a group of dislocations sufficiently large for crack formation; U is the activation energy for overcoming the dislocation barrier.

This relation for $W(T)$ goes through a sharp transition from 0 to 1 in a small temperature range near

$$T_c = \frac{U}{k \ln \frac{\nu l t_0}{b}}. \qquad (2)$$

The quantities in the logarithm can be estimated approximately; then, according to (2), for the experimentally obtained transition temperature $U \sim 1$ ev at the given deformation rate of 15% per minute. When these values are inserted into the expression for $W(T)$ we obtain the temperature dependence of the transition from the brittle to the plastic state, which is in good agreement with experiment (Fig. 4). Similar experiments with amalgamated single crystals of cadmium and lead did not reveal such a great reduction of strength. In these cases the lowering of the surface tension by mercury films, because of the low yield point P_f for these metals, did not satisfy the inequality $\sigma_{12} < \alpha b P_1$ under our experimental conditions.

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ON THE CLASSIFICATION OF SECONDARY EMISSION

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1. Through interactions with atoms and molecules a primary light beam is transformed into secondary emission which to some extent reveals properties of the substance under investigation. In some instances (complex molecules in solution, uranyls, etc.) the characteristics of the secondary emission are completely determined by the properties of the substance which transforms the light; in other instances (atoms, pairs of simple molecules) the character of the secondary emission depends on the incident radiation.

According to its observed characteristics all of the transformed light is divided into photoluminescence and Rayleigh and combination scattering. According to Vavilov [1] luminescence is emission whose duration is greater than that of the period of the light vibrations so that it can be quenched. On the other hand, scattering disappears (or appears) practically instantaneously after cessation (or beginning) of the action of the exciting force. Accordingly it is also assumed that scattering is not quenched.

Vavilov's criterion, the capacity for quenching, has enabled him to correctly qualify nonsteady emission [1]. This served as an experimental confirmation of his criterion. But the literature contains no theoretical basis for the criterion. No discussion can be found as to how thoroughly this criterion reflects the essence of light transformation processes.

2. In the classical theory the transformation of light is studied by analyzing the solution of the dipole equation [2]

$$\ddot{x} + 2(\gamma_e + \gamma_i)\dot{x} + \omega_0^2 x = \frac{e}{m} E_0 \cos \omega t. \tag{1}$$

Here ω_0 is the natural frequency of the dipole, γ_e and γ_i are damping constants for radiative damping and interaction with the medium. If at the instant when an external field is switched on the dipole is at rest, then, independently of the frequency of the incident light, steady emission cannot be established immediately but only after the time $t > 1/(2\gamma_e + 2\gamma_i)$. In a steady regime the dipole transforms light without changing the frequency according to

$$I(\omega) \sim \frac{\omega^4 I_0(\omega)}{(\omega_0^2 - \omega^2)^2 + 4(\gamma_e + \gamma_i)^2 \omega^2}, \tag{2}$$

where $I_0(\omega)$ and $I(\omega)$ are the intensities of the primary and of the secondary flux. Hence it follows that the position and shape of the transformed line is strongly dependent on the spectral form of the incident line. If irradiation is in the form of a narrow line ($\Delta\omega \ll \gamma_e + \gamma_i$) the secondary line coincides with the primary line in both frequency and shape ($I(\omega) \sim I_0(\omega)$). The ratio of the emitted to the absorbed energy determines the energy yield of the dipole emission [2]:

$$B = \frac{\gamma_e}{\gamma_e + \gamma_i}. \tag{3}$$

With increasing values of γ_i the yield declines. However, this affects the intensity of emission only near resonance; for $|\omega_0 - \omega| \gg \gamma_e + \gamma_i$, according to (2) the secondary intensity is independent of γ_i, that is, of the quencher.

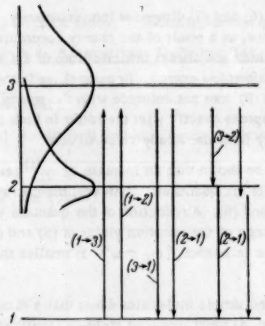

Fig. 1.

When irradiation is terminated the steady emission described by (2) disappears instantaneously ($I_0(\omega) = 0$) and the afterglow is observed with the duration $\tau = 1/(2\gamma_e + 2\gamma_i)$ and spectral form

$$S(\omega) \sim \frac{1}{(\omega_0 - \omega)^2 + (\gamma_{e\,i} + \gamma_i)^2}. \qquad (4)$$

Here $S(\omega)$ is the energy of frequency ω emitted by the dipole during the entire period of afterglow. It follows from (2) and (4) that in the general case the steady spectrum does not coincide with the afterglow spectrum. Only when the irradiating line is broad ($\Delta\omega \gg \gamma_e + \gamma_i$) and the frequency ω_0 is within that line do both spectra have the same shape. But even in this special case the transformation of light proceeds essentially in the same way as in the other cases.

3. The results of the calculation for the classical dipole also hold true in quantum electrodynamics. In addition, quantum electrodynamics enables us to investigate easily other possible transformations of light. For example, let us consider an atom possessing three energy levels and interacting with its ambient. The light transformed by this atom consists of three lines in a steady regime [3]. The arrows in Fig. 1 denote the most probable processes. The length of an arrow represents the frequency of the light involved in the process. The directions of the arrows and the symbols near them denote the atomic transitions; the absorption of a quantum of a given frequency can be associated with a transition to a second and to a third state. One of the secondary lines results from a transformation without frequency change: $\omega'' = \omega'$, where ω' is the incident frequency. The energy distribution in this line is described by

$$I_p(\omega'') \sim \left| \sum_{j=2,3} \left[\frac{d_{1j}^{(2)} d_{j1}^{(1)}}{-(\gamma_j^{(e)} + \gamma_j^{(\tau p)}) + (\omega_{j1} - \omega'')} + \frac{d_{1j}^{(1)} d_{j1}^{(2)}}{\omega_{j1} + \omega''} \right] \right|^2 \omega''^4 I_0(\omega''), \qquad (5)$$

where $d_{j1}^{(1)}$ and $d_{j1}^{(2)}$ are the matrix elements of the projection of the dipole moment on the polarization direction of the incident and secondary quanta, and $\gamma_j^{(e)}$ and $\gamma_j^{(i)}$ are the radiative and interaction damping constants of the jth level.

This line has characteristics almost identical with those of the classical line (2), except for the fact that the latter is obtained for a two-level system.

The second line gives the transformation when the frequencies of the secondary and primary quanta differ by the natural frequency of the atom: $\omega'' = \omega' - \omega_{21}$. When level 2 is narrow (for a metastable level) the intensity of this line is given by

$$I_k(\omega'') \sim \left| \frac{d_{23}^{(2)} d_{31}^{(1)}}{-i(\gamma_3^{(e)} + \gamma_3^{(i)}) + (\omega_{32} - \omega'')} + \frac{d_{23}^{(1)} d_{31}^{(2)}}{\omega_{21} + \omega''} \right|^2 \omega''^4 I_0(\omega'' + \omega_{21}). \qquad (6)$$

Equations (5) and (6) show that the positions and shapes of both lines depend strongly on the frequency and shape of the incident line. When the incident line is narrow ($\Delta\omega \ll \gamma_3^{(e)} + \gamma_3^{(i)}$) they agree in shape with the primary line independently of the degree of resonance ($\omega' \simeq \omega_{21}$).

The third line ($2 \rightarrow 1$) differs essentially from the first two. Its frequency and shape do not depend at all on the character of the incident line and are determined only by the properties of the atom. The intensity distribution in this line is given by

$$I_1(\omega'') = \frac{D}{(\omega_2 - \omega'')^2 + (\gamma_2^{(e)} + \gamma_2^{(i)})^2}, \qquad (7)$$

where D is the integral intensity of the line to within a factor (in general, it depends on the conditions of excitation). This is an intrinsic line of the atom take out power.

After the cessation of irradiation all of the steady lines (5), (6) and (7) disappear instantaneously, or, more precisely, as fast as the primary light is cut off. In their place, as a result of the energy accumulated by the atom during the steady regime, there appear afterglow lines which are always intrinsic lines of the atom. Each steady line can be associated with a definite fraction of the afterglow energy. In general, as in the classical theory, the afterglow spectrum corresponding to lines (5) and (6) does not coincide with the steady spectrum. On the other hand, the afterglow corresponding to line (7) agrees exactly with the latter in both frequency and shape. The integral intensity of (7) decreases exponentially from the steady value to zero.

By determining the quantum yields of (5), (6) and (7) it can be shown that an increase of $\gamma_j^{(i)}$ results in a reduction of the quantum yields, that is, all of the secondary lines are quenched. However, there is some difference between the quenching of (7) and the quenching of (5) and (6). A reduction of the quantum yield of (7) is always associated with a reduction of its intensity. A change of the quantum yields of (5) and (6), just as in the classical theory, is reflected in their intensity only for resonance ($|\omega_{31} - \omega'|$ is smaller than or comparable with $\gamma_3^{(e)} + \gamma_3^{(i)}$).

4. A theoretical analysis of light transformation by atoms and simple molecules shows that a classification according to quenching is impossible if by quenching we mean, as usual, reduced yield. Classification according to duration of afterglow is ambiguous and can sometimes lead to incorrect results. If this criterion is taken literally all secondary emission must be assigned to photoluminescence since all steady radiation is associated with definite afterglow. To be sure, the afterglow spectrum does not always coincide with or "fit" the steady spectrum. The "fit" of the spectra can be used in classifying light transformations.

As already mentioned, line (7) differs essentially from lines (5) and (6). The frequency and shape of (7) are determined only by the properties of the atom so that the corresponding afterglow line fits well. On the other hand the shapes of (5) and (6) depend strongly on the frequency and shape of the incident line and thus differ generally from the afterglow spectrum. This difference is an external expression of the fact that these lines result from a different process mechanism. In the formation of lines (5) and (6), between the acts of primary absorption and secondary emission there are no intermediate processes. Therefore the two acts are not independent of each other; this is reflected in the close relation between the primary and secondary quanta. On the other hand, the emission of the quanta in (7) is separated from the act of absorption by the transition $3 \rightarrow 2$, which accompanies the emission of quanta with the combination frequency $\omega' - \omega_{31}$. This intermediate process eliminates all dependence of the secondary light on the character of the primary light.

In view of the above the classification of light transformations can proceed as follows. Line (7) must be assigned to photoluminescence, but lines (5) and (6) must be assigned to Rayleigh and Raman scattering. In nonresonant irradiation the emission described by (5) and (6) is certainly scattering and in resonance becomes more intense. In this case line (5) is usually called "resonance radiation" or "resonance fluorescence" and line (6) is called "fluorescence" [4, 5]. We believe that the transformation of light described by (5) and (6) should in the case of resonance be called resonance scattering since its mechanism does not differ essentially from that of light transformation in the nonresonant case. There are only superficial differences between the two cases. The transition of resonance radiation to specular reflection with increased vapor pressure [4, 5] supports this classification.

The properties of light transformed by a classical dipole are analogous to those of the radiation described by (5). Therefore the classical dipole enables us to describe only Rayleigh scattering.

5. Through interactions with the ambient the atoms and simple molecules can change their state without losing the ability to emit light. The intermediate processes lead to quenching of the scattering and to a correspondingly increased intensity of some photoluminescence lines. In complex molecules there is also a rapid redistribution of the vibrational and rotational energy of the molecule over its degrees of freedom [2]. Calculations show that as a result the scattering is completely quenched so that only photoluminescence is ever observed.

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NEGATIVE K-PARTICLE CAPTURE WITH THE EMISSION OF $_{\Lambda}He_2^5$

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(Presented by Academician N. N. Bogoliubov June 1, 1957)

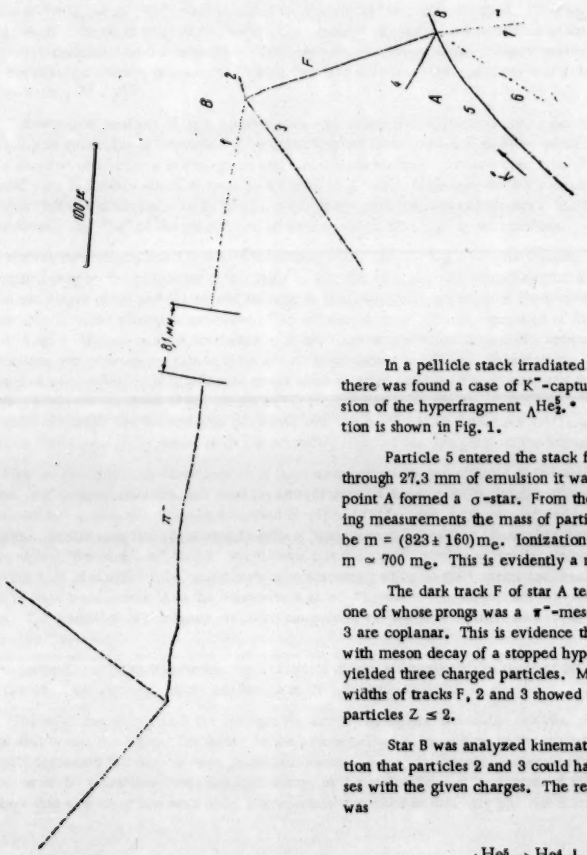

Fig. 1.

In a pellicle stack irradiated at a high altitude there was found a case of K^--capture followed by emission of the hyperfragment $_{\Lambda}He_2^5$.* The microprojection is shown in Fig. 1.

Particle 5 entered the stack from below. Passing through 27.3 mm of emulsion it was stopped and at point A formed a σ-star. From the range and scattering measurements the mass of particle 5 was found to be $m = (823 \pm 160)m_e$. Ionization measurements gave $m \simeq 700\ m_e$. This is evidently a negative K-particle.

The dark track F of star A terminates in star B, one of whose prongs was a π^--meson. Prongs 1, 2, and 3 are coplanar. This is evidence that star B was formed with meson decay of a stopped hyperfragment which yielded three charged particles. Measurements of the widths of tracks F, 2 and 3 showed that for each of these particles $Z \leq 2$.

Star B was analyzed kinematically on the assumption that particles 2 and 3 could have all possible masses with the given charges. The resulting decay scheme was

$$_{\Lambda}He_2^5 \rightarrow He_2^4 + p + \pi^-.$$

* Discovered by N. V. Kirsanova.

TABLE 1

Track	F	1	2	3	4	5	6	7	8
True range, μ	161	8481	19.3	401	54.7	>27300	>27500 unstopped	4880	18
Error of range measurement, %	0.4	0.1	4.2	0.5	1.7	—	—	0.1	5.1
Spread of ranges, %	1.2	3.0	2.0	1.6	1.8	—	—	1.4	2
Angle β with horizontal plane in unprocessed emulsion, deg.*	0	−9.0	+12.4	−18.5	+16.0	—	+7.9	−19.8	+15.8
Error of measurement of β, deg.	—	4	7	2	2	—	0.3	1	7
Azimuthal angle φ, deg.**	0	232,3	79.5	278.2	44.5	—	126.6	167.8	269,8
Error of measurement of φ, deg.	0.7	0.7	1.7	0.7	0.7	—	0.7	0.7	1.7
Number of δ electrons with range > 4 grains***	0	—	0	0	0	—	—	—	0
Charge Z	$\leqslant 2$	—	$\leqslant 2$	$\leqslant 2$	$\leqslant 2$	−1	±1	+1	$\leqslant 2$
Method of measuring mass and its value, m_e	—	—	—	—	—	$\langle\alpha\rangle - R$: 914 ± 170; $I - R$: 658 ± 55	$I - \langle\alpha\rangle$: 343^{+94}_{-62}	$\langle\alpha\rangle - R$: 1250 ± 340; $I - R$: 2700 ± 300	—
Identification of particle	${}_{\Lambda}He_2^5$	π^-	He_2^4	p	p	K^-	π^-	p	He_2^4
Energy of particle, Mev	21.3+0.3	21.0+0.4	4.8±0.2	8.4±0.1	2.47±0.04	—	66.3+6.5	35.3+0.3	4.6+0.2

* The sign "+" is used for tracks directed toward the surface; the sign "−" is used for tracks toward the glass.

** With respect to an arbitrarily chosen direction, which is identical for all star prongs and the track of the hypernucleus.

*** 1 grain in the track, 3 grains outside the track.

The kinetic energy of the decay products was $Q_k = (34.2 \pm 0.4)$ Mev. The sum of the momenta of the produced particles was $p = (13 \pm 26)$ Mev/c. The binding energy was $B_\Lambda = (2.7 \pm 0.4)$ Mev, which is in good agreement with published data. The angle between the normal to the plane of disintegration and the direction of flight of the hyperfragment was 82°.

In this case as in other published cases of ${}_\Lambda He_2^5$ it can, of course, be postulated that a neutron with very little energy was also emitted. This would not affect the coplanarity of the prongs of star B. The decay scheme would then be

$$ {}_\Lambda He_2^6 \rightarrow He_2^4 + p + n + \pi^-. $$

However, the kinematical analysis of primary star B which is carried out below enables us to make a more exact identification of hyperfragment F.

If particle 8 were a proton or alpha particle its energy would be 1.19 or 4.6 Mev, respectively, which is considerably smaller than the Coulomb potential barrier in the heavy atoms of the emulsion. It was therefore assumed that the neutron which captured the K^--meson was C^{12}, N^{14} or O^{16}.

For all possible (according to Table 1) values of the masses and charges of particles 4, 6, 7 and 8 and for both types of hyperfragments (${}_\Lambda He_2^5$ and ${}_\Lambda He_2^6$) we considered all variants of the reactions involving C^{12}, N^{14} and O^{16}. Momentum and energy balance showed that the conservation laws are satisfied only by

$$ K^- + C_6^{12} \rightarrow {}_\Lambda He_2^5 + p + \pi^- + p + He_2^4 + n \qquad (1) $$

and possibly

$$ K^- + C_6^{12} \rightarrow {}_\Lambda He_2^5 + p + \pi^- + p + He_2^3 + 2n. \qquad (2) $$

(1) and (2) differ in the number of emitted neutrons and in the type of particle 8. In both instances F is ${}_\Lambda He_2^5$.

From Reaction (1) we obtain for K^- the mass $m_{K^-} = (494.3 \pm 6.8)$ Mev, which is in good agreement with data in other papers.*

From (2) it was possible to obtain a lower limit of the K^- mass; this value was (506 ± 6.6) Mev. The lower limit was calculated on the unlikely assumption that both neutrons were emitted in the same direction with equal momenta. Otherwise the calculated value of m_{K^-} is still greater. For this reason Reaction (2) is excluded as unlikely.

Without conflicting with current theoretical notions it can be assumed that the formation of a hyperfragment can sometimes involve a Σ-particle.** A Σ-particle formed within a nucleus can emerge as part of the hyperfragment. In a very short time it is transformed into a Λ^0-particle, as a result of which the initial hyperfragment decays into a lighter hyperfragment containing a Λ^0-particle and one or more nucleons.***

If in our case the hyperfragment was produced in this manner the track of the ${}_\Lambda He_2^5$ and one or more prongs of star A were produced by particles in the decay of the original hyperfragment with a Σ-particle. The momentum of these particles must equal the momentum of the hyperfragment with the Σ-particle before decay. The energy of the hyperfragment can be determined from its momentum. The total energy of the original hyperfragment is equal to the total energy of the particles emitted in its decay. This enables us to determine B_Σ, the binding energy of the Σ-particle.

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* $m_{K^+} = (493.66 \pm 0.36)$ Mev [1]. The mass of K^--mesons is known with less accuracy. However, it is reasonable to assume $m_{K^-} = m_{K^+}$.

** The energy of a π-meson formed in a σ_K-star is quite small (53 Mev). This does not contradict the hypothesis that in the K-particle capture of a Σ-particle was first formed.

*** This possibility is meaningful only if in the time required for the transition from a Σ-particle to a Λ^0-particle the hyperfragment is able to escape from the original nucleus. This is not excluded in principle by the Gell-Mann scheme.

In this manner we tested all possible combinations of 2, 3, 4 and 5 particles of the primary star. For almost all combinations B_Σ assumed either negative or very high positive values; for only two combinations were acceptable values obtained (of the order of a few Mev).

1st combination. ${}_\Lambda He_2^5$, n and 2p were formed in the decay of the hyperfragment ${}_\Sigma Be_4^8$.

$B_\Sigma = (5.9 \pm 2.0)$ Mev if the hyperfragment included Σ^0.

$B_\Sigma = (5.5 \pm 2.0)$ Mev if the hyperfragment included Σ^+.

2nd combination. ${}_\Lambda He_2^5$, p, n and He_2^4 were formed in the decay of ${}_\Sigma B_5^{11}$.

$B_\Sigma = (1.3 \pm 2.0)$ Mev if the hyperfragment included Σ^0.

$B_\Sigma = (0.2 \pm 2.0)$ Mev if the hyperfragment included Σ^+.

The result does not, of course, mean that the hyperfragment was necessarily formed in the manner described above, since the values of B_Σ may have been small only by chance. This question can be answered only by a similar analysis of other cases. It would be especially useful to examine σ_K-stars which permit a complete kinematical analysis [2]. It is also possible that a hyperfragment is sometimes formed in an excited state with subsequent transition to the ground state and emission of a gamma ray or nucleon.* Using an analysis similar to the above and with sufficient material it would be possible to attempt to detect such excited hyperfragments.

In our specific case there are two combinations of particles for which the binding energy B_Λ^* becomes negative.

1st combination. ${}_\Lambda He_2^5$ and p were produced in the decay of the excited hyperfragment ${}_\Lambda Li_3^{*6}$. Binding energy $B_\Lambda^* = (2.2 \pm 0.7)$ Mev. Proton energy in the center-of-mass system $E_p = (10.6 \pm 0.2)$ Mev.

2nd combination. ${}_\Lambda He_2^5$ and n were produced in the decay of the excited hyperfragment ${}_\Lambda He_2^{*6}$. Binding energy $B_\Lambda^* = (-0.9 \pm 2.0)$ Mev. Neutron energy in the center-of-mass system $E_n = (9.9 \pm 1.1)$ Mev.

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* As when the compound nucleus Li_3^5 becomes He_2^4 with emission of a proton [3] or He_2^5 becomes He_2^4 with emission of a neutron [4].

DISPERSION RELATIONS FOR THE SCATTERING OF K-MESONS BY NUCLEONS

M. K. Polivanov

(Presented by Academician N. N. Bogoliubov May 15, 1957)

1. In the investigation of strong interactions the dispersion relations provide at present the only means of obtaining positive even though incomplete information; the assumptions which are required to obtain these relations amount essentially to admission of the existence of a S matrix with the usual general properties and to the strict observance of causality (local interactions). We shall accordingly analyze here the dispersion relations for the scattering of K-mesons by nucleons, $K + N \to K + N$ (and $\tilde{K} + N \to \tilde{K} + N$), without mentioning familiar details.

2. Considering the K-meson to be a scalar particle (θ-meson) or pseudoscalar particle (τ-meson), we associate with it the corresponding field $\kappa_\rho(x)$. The subscript ρ denotes the four states which are associated linearly with the four particles (K^+, K^0, $\tilde{K}^0$, K^-) differing in their isotopic spin projection and hypercharge (or strangeness). In isotopic space these four fields form a double spinor like the four functions of a nucleon and antinucleon. We shall consider the Yukawa interaction of K-mesons and nucleons: $N \leftrightarrow K + Y$, $Y \leftrightarrow N + \tilde{K}$. It will become clear below that other possible interactions are not explicitly reflected in the dispersion relations.

The amplitude of the scattering process $K + N \to K + N$ can be represented in a familiar manner by

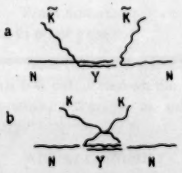

Fig. 1.

$$S_{\alpha\omega}(q, q') = \langle \mathbf{p}'s', \mathbf{q}'\rho' | S | \mathbf{p}s, \mathbf{q}\rho \rangle =$$

$$= \frac{1}{(2\pi)^3 \sqrt{4q^0 q'^0}} \int e^{i(q'x - qy)} \langle \mathbf{p}'s' \left| \frac{\delta^2 S}{\delta \kappa_\rho(x)\, \delta \kappa_{\rho'}(y)} S^+ \right| \mathbf{p}s \rangle\, dx\, dy, \qquad (1)$$

$$p^0 = \sqrt{\mathbf{p}^2 + M^2}, \ldots, q'^0 = \sqrt{\mathbf{q}^2 + m^2},$$

where $\mathbf{p}$, $\underline{s}$, $\mathbf{q}$ and ρ are the momentum and other quantum numbers of the nucleon and meson before scattering, and the same symbols when primed denote the quantum numbers after scattering; α and ω are combined subscripts of the initial and final states.

Using the equality

$$\frac{\delta^2 S}{\delta \kappa_\rho(x)\, \delta \kappa_{\rho'}(y)} S^+ = -\, i_\rho(x)\, i_{\rho'}(y) - i \frac{\delta i_\rho(x)}{\delta \kappa_{\rho'}(y)}, \qquad (2)$$

where $i_\rho(x) = i \dfrac{\delta S}{\delta \kappa_\rho(x)} S^+$ is the current operator, we find that in the region of real momenta $\mathbf{p}_i^2 > 0$ and positive energies $p_i^0 > 0$ associated with the masses by the usual relation $p_i^2 = m_i^2$, the amplitude $S_{\alpha\omega}$ agrees with

$$H_{\alpha\omega}(q, q') = \frac{-i}{(2\pi)^3 \sqrt{4q_0 q'_0}} \int dx\, dy\, e^{i(q'x - qy)} \langle \mathbf{p}'s' \left| \frac{\delta i_{\rho'}(y)}{\delta \kappa_\rho(x)} \right| \mathbf{p}s \rangle. \qquad (3)$$

From translational invariance this becomes

$$H_{\alpha\omega}(q, q') = \frac{2\pi i}{\sqrt{4q^0 q'^0}}\,\delta(p + q - p' - q')\,T^{ret}_{\alpha\omega}\left(\frac{q+q'}{2}\right). \tag{4}$$

Introducing the system of coordinates where $\mathbf{p} + \mathbf{p}' = 0$, we have

$$\begin{gathered} \mathbf{q} = \mathbf{p} + \lambda\mathbf{e}, \quad \mathbf{q}' = -\mathbf{p} + \lambda\mathbf{e} \quad (\mathbf{e}^2 = 1, (\mathbf{e}\cdot\mathbf{p}) = 0), \\ q^0 = q'^0 = \sqrt{m^2 + \mathbf{p}^2 + \lambda^2} = E, \quad p^0 = p'^0 = \sqrt{\mathbf{p}^2 + M^2}, \quad \lambda = \sqrt{E^2 - m^2 - \mathbf{p}^2}, \end{gathered} \tag{5}$$

where T^{ret} is a function of E and $\lambda\mathbf{e}$.

Using the symmetrization $\mathfrak{S}_+ f = f(\mathbf{e}) + f(-\mathbf{e})$, $\mathfrak{S}_- f = \frac{1}{\lambda}\{f(\mathbf{e}) - f(-\mathbf{e})\}$ to eliminate the ambiguity introduced by λ, we can continue $\mathfrak{S}T^{ret}$ analytically in the region of complex values of E.

Bogoliubov's rigorous proof of analyticity for π-mesons [1] must be specially modified for the present case; however, we shall not do this here, but shall postulate that the function $\mathfrak{S}T^{ret}(E)$ obtained by analytic continuation possesses the properties which enable us to describe the dispersion relations in the usual form by representing $\mathrm{Re}\,T^{ret}_{\alpha\omega} = D_{\alpha\omega}$ through the integral of $\mathrm{Im}\,T^{ret}_{\alpha\omega} = A_{\alpha\omega}$ along the entire real axis of E.

In order to obtain relations which relate only physically observable quantities we must reduce the integration interval. Forming the combinations $(1 \pm P_{\rho\rho'})\,\mathfrak{S}\,T^{ret}_{\alpha\omega}$ (where $P_{\rho\rho'}$ is the permutation operator of the subscripts ρ and ρ'), which have definite parity with respect to E, we can eliminate integration over negative energies. Finally, by fixing the degree of increase of the amplitude for $E \to \infty$, assuming a linear increase,* we obtain four dispersion relations for the amplitudes with different symmetries, one of which is given here as an illustration:

$$(1 + P_{\rho\rho'})\,\mathfrak{S}_+\{D_{\alpha\omega}(E) - D_{\alpha\omega}(E_0)\} = \frac{2}{\pi}(E^2 - E_0^2)P\int_0^\infty dE\,\frac{E'\mathfrak{S}_+(1+P_{\rho\rho'})\,A_{\alpha\omega}(E')}{(E'^2 - E^2)(E'^2 - E_0^2)}.$$

3. However, even after eliminating negative energies the integration interval contains the region $0 < E < \sqrt{m^2 + \mathbf{p}^2}$ which does not pertain to scattering states and which can thus not be determined from a scattering experiment. For the analysis of this region we use a representation of $A_{\alpha\omega}(E)$ which follows from (2) and the and the hypothesis that a complete set of states exists:

$$\begin{aligned} A_{\alpha\omega}(E, \mathbf{e}) = (2\pi)^4 i \sum_n \langle \mathbf{p}'s' | i_\rho(0) | \lambda\mathbf{e}, n\rangle \langle \lambda\mathbf{e}, n | i_{\rho'}(0) | \mathbf{p}s\rangle \cdot \\ \delta(\sqrt{M^2 + \mathbf{p}^2} + E - \sqrt{M_n^2 + \lambda^2}) - (2\pi)^4 i \sum_n \langle | \mathbf{p}s' | i_{\rho'}(0) | | -\lambda\mathbf{e}, n\rangle \cdot \\ \cdot \langle -\lambda\mathbf{e}, n | i_\rho(0) | \mathbf{p}s\rangle \delta(-\sqrt{M^2 + \mathbf{p}^2} + E + \sqrt{M_n^2 + \lambda^2}). \end{aligned}$$

The behavior of $A_{\alpha\omega}(E)$ is determined by the δ-functions of the energy which it contains. The arguments of the δ-functions vanish at the points

$$\pm E = \frac{M_n^2 - M^2 - m^2 - 2\mathbf{p}^2}{2\sqrt{M^2 + \mathbf{p}^2}}.$$

* This assumption, for which the total cross section approaches a constant, is very reasonable. It is the assumption for the scattering of pions and can apparently be regarded as common to all strong interactions. It is not a limiting assumption and we can allow for any degree in accordance with experiment.

The mass spectrum of possible intermediate states can thus be determined. Conservation of strangeness and nucleonic charge allow the states: hyperon Y (Λ or Σ), hyperon and pion (Y + π) , nucleon and K-meson (N + K),* etc.

The Y state belongs to the discrete spectrum; its contribution is easily calculated and it is proportional to the square of the NKY interaction constant. Thus, for example, the symbolically written contribution of the Λ-hyperon for pseudoscalar coupling is

$$\Gamma(\Lambda) = G^2 \left\{ \lambda_{s's} \frac{2\mathfrak{M}(\mathbf{p}^2 + M^2) - (\mathfrak{M}^2 + M^2 - m^2) M}{4(M^2 + \mathbf{p}^2)^{3/2}} - \right.$$

$$\left. - i\lambda_p \frac{(\vec{\sigma} \cdot [\mathbf{p} \times \mathbf{e}])\, s's}{2(M^2 + \mathbf{p}^2)} \right\} \delta(E - E_p),$$

but the continuous spectrum begins with the lowest two-particle state. Since the limit of the continuous spectrum

$$E_p^{\Lambda+\pi} = \frac{(\mathfrak{M} + \mu)^2 - M^2 - m^2 - 2\mathbf{p}^2}{2\sqrt{M^2 + \mathbf{p}^2}}$$

lies below the limit of the scattering states $E_1 = \sqrt{m^2 + \mathbf{p}^2}$ even for the most favorable case of forward scattering (i. e., $\mathbf{p} = 0$ in our system), it is still necessary to obtain some estimate of the integral from $E_p^{\Lambda+\pi}$ to E_1. For this purpose we can use an analytic expression for $A_{\alpha\omega}$ in the low-energy region ($E > m$, $E \sim m$) and then regard this expression as valid in the region of purely imaginary momenta ($E < m$). Without discussing the validity of this procedure we can expect it to be suitable at least for a rough estimate.

4. If we postulate any other interactions such as the direct interaction of pions and K-mesons according to Goldhaber and Schwinger's hypothesis [2] it is evident that these must not be taken explicitly into the dispersion relations, because the lowest state for such an interaction will lie above the threshold for scattering states and will be taken into account automatically in the experimental $A_{\alpha\omega}(E)$ for $E > m$.

Thus the contribution from the hyperon band enables us to evaluate the YKN interaction constant. Unfortunately, this result is somewhat smeared because of the existence of an indefinite contribution from the continuous spectrum.

I am deeply grateful to Academician N. N. Bogoliubov for suggesting this problem and for his continued interest, and to my associates for discussions and assistance.

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* The first two states (Y and Y + π) are allowed only in the process $\bar{K} + N \rightarrow \bar{K} + N$.

ON THE INTERACTIONS OF MUONS AND NUCLEONS

I. S. Shapiro, E. I. Dolinskii and L. D. Blokhintsev

(Presented by Academician D. V. Skobeltsyn May 27, 1957)

Concerning nonelectromagnetic interactions of muons and nucleons we at present have very little experimental data, from which there follows only the fact that this interaction is "weak." Yet a more careful study of the muon-nucleon interaction is of great importance for understanding the nature of the muon, which is one of the most puzzling elementary particles that we are acquainted with (the muon and electron are characterized by the same types of interaction – weak and electromagnetic; both particles possess the same spin but the muon mass is 206.7 times that of the electron). In particular, the clarification of the form of muon-nucleon interactions may be decisive in explaining why, although pions and K-mesons decay to μ and ν,* these particles do not decay to $\underline{e}$ and ν. Indeed, if pion and K-meson decays go through a baryon vacuum, then, regarding pions and K-mesons as pseudoscalar particles, it is easily shown that decay is forbidden for scalar, vector and tensor four-fermion interactions but allowed for the pseudoscalar and pseudovector variants. Since only the scalar and tensor variants are apparently responsible for electron-nucleon interactions,** as follows from beta decay, the decay of pions and K-mesons to $\underline{e}$ and ν will be forbidden. The presence of decays to μ and ν would be evidence for pseudoscalar and pseudovector interactions between muons and nucleons (or at least one of these types).

If, on the other hand, pion and K-meson decay occurs even partially through a direct interaction between the pion or K-meson with a muon and neutrino, the muon-nucleon interaction Hamiltonian must include a pseudoscalar interaction.

It is well known that when muons are slowed down by matter they are absorbed by nuclei and that this process passes through an intermediate stage with the formation of a mesic atom. It is customarily assumed that the absorption of negative muons takes place through the reaction

$$\mu^- + p \to n + \nu. \tag{1}$$

In the present paper we consider the angular distribution of neutrons resulting from negative muon capture by the proton in μ-mesic hydrogen. It is assumed that the μ^--meson is polarized.*** In this case, by virtue of parity nonconservation in weak interactions, the angular distribution of neutrons will generally be anisotropic and the sign and magnitude of the anisotropy will depend on the type of interaction.

The muon-nucleon interaction energy with parity nonconservation can be written as

$$H = \sum_k (\bar{\psi}_n O_k \psi_p)\,(\bar{\psi}_\nu [g_k - g_k' \gamma_5]\, O_{ki} \psi_\mu) + \text{H.C.} \tag{2}$$

---

* The symbol ν denotes a neutrino.

** At present in the light of recent experimental data the selection of the types of interaction in beta decay has become more complicated.

*** The negative mesons which are created in negative pion decay are polarized parallel to their direction of motion [1]. With slowing down and the formation of a μ-mesic atom depolarization is possible, but not complete depolarization [2]. Of course, an experiment on the angular distribution in Reaction (1) should be preceded by a measurement of the polarization of the negative muon in the mesic atom (for example, by measuring the anisotropy in the angular distribution of decay electrons).

where the O_k are operators known from beta decay theory, which are composed of the Dirac matrices γ_k; k = s, p, v, a, t, where s, p, v, a, t denote, respectively, scalar, pseudoscalar, vector, pseudovector and tensor interactions. When $g_k = -g_k$ we obtain Landau's version of the theory [1] with a longitudinally polarized neutrino.

The angular distribution of neutrons is described by

$$W(\theta) = 1 + \alpha \cos\theta, \tag{3}$$

where θ is the angle between the direction of flight of the neutron and the direction of polarization of the negative muon.

In the general case when all types of interaction are present α has the form

$$\alpha = A/B, \tag{4}$$

where

$$\begin{aligned} A = {} & (E+M)h_{ss} + (E-M)h_{pp} + 2(M+\varepsilon)h_{vv} - 2(M-\varepsilon)h_{aa} - \\ & - 8(E+2\varepsilon)h_{tt} + 2(E+M+\varepsilon)\operatorname{Re}h_{sv} + 4\varepsilon\operatorname{Re}h_{st} + 2(E-M+\varepsilon)\operatorname{Re}h_{ap} + \\ & + 4\varepsilon\operatorname{Re}h_{tp} - 4\varepsilon\operatorname{Re}h_{va} - 4(E-M+\varepsilon)\operatorname{Re}h_{vt} - 4(E+M+\varepsilon)\operatorname{Re}h_{at}, \end{aligned} \tag{4a}$$

$$\begin{aligned} B = {} & (E+M)f_{ss} + (E-M)f_{pp} + 2(2E-M+\varepsilon)f_{vv} + \\ & + 2(2E+M+\varepsilon)f_{aa} + 8(3E+2\varepsilon)f_{tt} + 2(E+M+\varepsilon)\operatorname{Re}f_{sv} + \\ & + 4\varepsilon\operatorname{Re}f_{st} + 2(E-M+\varepsilon)\operatorname{Re}f_{ap} + 4\varepsilon\operatorname{Re}f_{tp} + 4\varepsilon\operatorname{Re}f_{va} + \\ & + 12(E-M+\varepsilon)\operatorname{Re}f_{vt} + 12(E+M+\varepsilon)\operatorname{Re}f_{at}. \end{aligned} \tag{4b}$$

Here

$$f_{ik} = g_i^* g_k + g_i'^* g_k', \quad h_{ik} = g_i^* g_k' + g_i'^* g_k; \tag{5}$$

E, M are the total energy and rest energy of the neutron, respectively; ϵ is the neutrino energy (E = 944.7 Mev, M = 939.5 Mev, ϵ = 99.1 Mev).

The coefficient α for the different types of interaction (2) with a longitudinal neutrino is given in Table 1.*

TABLE 1

Type of interaction	s	p	v	a	t
α	−1	−1	−0.99	+0.29	+0.38
P	0	0	0.095	0.64	0.72

Because of the special importance of the combinations s + t and p + a we give the values of α for these types separately. If in accordance with the hypothesis of the conservation of combined parity, we consider the scalar and tensor interaction constants as real and use their values from beta decay of nucleons (see [4], for example): $g_s = (+1.4 \text{ or } -1.4)\cdot 10^{-49}$ erg-cm^3, $g_t = +1.8\cdot 10^{-49}$ erg-cm^3, then for a longitudinal neutrino we obtain $\alpha^{s+t} = +0.30$ or $+0.32$. Similarly in the case of a longitudinal neutrino for the combination p + a we obtain

$$\alpha^{p+a} = \frac{1.78\,|g_a|^2 - 0.0055\,|g_p|^2 - 0.22\operatorname{Re} g_a^* g_p}{6.20\,|g_a|^2 + 0.0055\,|g_p|^2 + 0.22\operatorname{Re} g_a^* g_p}.$$

* According to experiment [3] an antineutrino emitted in the β^--decay of nucleons is polarized counter to its direction of motion. Our form of the interaction Hamiltonian corresponds to emission in Process (1), just as for electron capture by nuclei, of a neutrino polarized in its direction of motion. When an antineutrino is emitted in μ^--capture the sign of α is reversed.

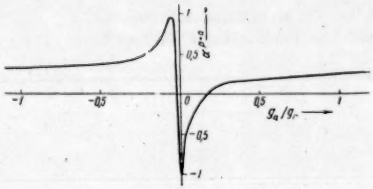

Fig. 1.

The values of α^{p+a} as a function of $x = g_a/g_p$ are given in Fig. 1 for real values of x. It must be noted that the values of α for comparable values of g_p and g_a will be very close to α^a since the pseudoscalar constant always contains a small factor v/c (v is the neutron velocity). On the other hand, if the experimental value of α is close to α^p it can be stated confidently that the admixture of pseudovector and tensor interactions is negligibly small.

We note that similar formulas can be obtained for μ^--capture by protons bound in nuclei. The coefficient α will then generally depend on nuclear matrix elements, thus complicating the interpretation of experimental data. For some light nuclei these matrix elements can be evaluated when the final state of the resultant nucleus is known after Process (1).

In addition to using the anisotropy of the angular distribution of neutrons, we can determine the type of interaction from the fact that the neutrons resulting from (1) are generally polarized. This polarization can be either transverse or longitudinal, and part of the latter results from parity nonconservation and can occur even if the captured negative muons are not polarized. Table 1 gives the values of the longitudinal polarization of P neutrons resulting from the capture of unpolarized negative muons by free protons in the case of a longitudinal neutrino.

These findings are approximately applicable to muon capture by nuclei since the nuclear spin-orbit interaction has little effect on the longitudinal polarization of the neutron and its depolarization in the Coulomb field of the resultant nucleus is negligibly small.

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SIGNIFICANCE OF ABBREVIATIONS MOST FREQUENTLY ENCOUNTERED IN SOVIET PHYSICS PERIODICALS

AN SSSR	*Academy of Sciences, USSR*
FIAN	*Physics Institute, Academy of Sciences USSR*
GITI	*State Scientific and Technical Press*
GITTL	*State Press for Technical and Theoretical Literature*
GOI	*State Optical Institute*
GONTI	*State United Scientific and Technical Press*
Gosenergoizdat	*State Power Press*
Gosfizkhimizdat	*State Physical Chemistry Press*
Gozkhimizdat	*State Chemistry Press*
GOST	*All-Union State Standard*
Goztekhizdat	*State Technical Press*
GTTI	*State Technical and Theoretical Press*
GUPIAE	*State Office for Utilization of Atomic Energy*
IF KhI	*Institute of Physical Chemistry Research*
IFP	*Institute of Physical Problems*
IL	*Foreign Literature Press*
IPF	*Institute of Applied Physics*
IPM	*Institute of Applied Mathematics*
IREA	*Institute of Chemical Reagents*
ISN (Izd. Sov. Nauk)	*Soviet Science Press*
IIaP	*Institute of Nuclear Studies*
Izd	*Press (publishing house)*
KISO	*Solar Research Commission*
LETI	*Leningrad Electrotechnical Institute*
LFTI	*Leningrad Institute of Physics and Technology*
LIM	*Leningrad Institute of Metals*
LITMiO	*Leningrad Institute of Precision Instruments and Optics*
Mashgiz	*State Scientific-Technical Press for Machine Construction Literature*
MATI	*Moscow Aviation Technology Institute*
MGU	*Moscow State University*
Metallurgizdat	*Metallurgy Press*
MOPI	*Moscow Regional Institute of Physics*
NIAFIZ	*Scientific Research Association for Physics*
NIFI	*Scientific Research Institute of Physics*
NIIMM	*Scientific Research Institute of Mathematics and Mechanics*
NII ZVUKSZAPIOI	*Scientific Research Institute of Sound Recording*
NIKFI	*Scientific Institute of Motion Picture Photography*
OIIaI	*Joint Institute of Nuclear Studies*
ONTI	*United Scientific and Technical Press*
OTI	*Division of Technical Information*
OTN	*Division of Technical Science*
RIAN	*Radium Institute, Academy of Sciences of the USSR*
SPB	*All-Union Special Planning Office*
Stroiizdat	*Construction Press*
URALFTI	*Ural Institute of Physics and Technology*

NOTE: Abbreviations not on this list and not explained in the translation have been transliterated, no further information about their significance being available to us.—*Publisher.*

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